

**PHENOMENOLOGICAL INVESTIGATIONS
OF MULTIPARTICLE PRODUCTION IN HADRONIC COLLISIONS**

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Abstract The aim of the thesis is to study how the ordinary <u>low $p_{\perp}$</u> multiparticle production in hadronic collisions is generated. We work with two different models. Variations of two of the basic assumptions in the Landau hydrodynamic model is investigated. It is found that this model is unstable against variations of the <u>initial state</u> . On the other hand, the results of the model are insensitive to the <u>velocity of sound</u> being temperature dependent or not. We also present a model for baryon fragmentation, where we explain the low $p_{\perp}$ jet observed in data as the result of a quark fragmentation mechanism. A Monte Carlo generation program that produces jets in accordance with the model is described, and we find that the <u>longitudinal particle distributions</u> thus obtained are in good agreement with the experimental data for both mesons and baryons.			
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MULTIPARTICLE PRODUCTION IN HADRONIC COLLISIONS

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Thesis for the degree of Doctor of Philosophy
(Avhandling för doktorsexamen)

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- I. Generalized Hydrodynamical Models
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- III. A Model for the Reaction Mechanism and the Baryon
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INTRODUCTION

One of the aims of particle physics is to understand what matter is and to reveal its finest structures. Of the known particles, some, as the electron, are pointlike, (i.e. structureless) down to very small distances ($\sim 10^{-18}$ m). On the other hand, the great majority of particles, the hadrons (i.e. those which are strongly interacting), are known to be extended objects with dimensions of about 1 fermi = 10^{-15} m.

The most well-known hadrons are those that build up the atomic nucleus, i.e. the proton, p , the neutron, n , and the π -meson, which transmits the strong interaction that holds the nucleus together. Now the discovery in the 50's of a large number of particles as elementary as the p , n and π led to the concept of quarks, the constituents of hadrons. Quarks were introduced for pure systematic reasons, i.e. they made it possible to group particles into families and gave some understanding of the way they interact. The postulation of these remarkable objects launched an intense experimental search for them. Being fractionally charged they would be easily detected if they existed. But every search thus far has been futile, and if free quarks exist, they are immensely rare.

At short distances though ($\lesssim 1$ fermi) we have acquired evidence for the existence of pointlike objects, which can be identified with quarks. This has been revealed by so-called deep inelastic lepton-proton scattering experiments. And so gradually a picture of hadrons has emerged, depicting them as bags inside of which quarks can move relatively freely but from which they cannot escape, i.e. they are confined to the bags. Understanding this confinement mechanism is one of the main problems facing particle physics today.

Now the only way we know of investigating the structures of particles is by letting them collide at very high energies. The common feature of these collisions is that a large number of secondary particles are produced, so-called multiparticle production. This is the case also for hadronic collisions, where the great majority of events are characterized by the production of many particles with limited transverse momenta. This means in a way that the secondary particles are coming out in the form of two jets along the directions of the incoming particles. Usually one views the collisions in the center of mass system (cms), and in this reference frame most of the particles produced have low momentum and belong to what is called the central region. This is characterized by having properties that do not depend on the beam or the target. The rest of the particles, those emitted with high momentum, are called beam and target fragments, respectively, if they move forward or backward in the cms. The reason for this is that they to some extent remember the quantum numbers of the beam and target particle respectively.

The most consistent way that we know to describe the creation and annihilation of matter is to use quantum field theory. The strong interaction has however for a long time evaded the attempts of physicists to do this. But today there exists one such theory called QCD, Quantum ChromoDynamics. In this theory the quarks are glued together by eight so-called gluons. Both quarks and gluons carry a special kind of charge called color charge, but the hadrons themselves are color neutral. QCD explains why quarks are free at distances within a hadron, and the hope is that it also will explain why they cannot get out of the hadron bag. So far, however, the mathematical difficulties of the theory have been insurmountable and its predictiveness therefore to a large extent limited.

Over the years, the difficulties to achieve a profound theoretical description of strong interactions have all along made the use of

phenomenological models a necessity. The models teach us more about the interactions and help us to figure out what dynamical mechanisms are responsible for the outcome of different types of experiments that we carry out. So although the understanding of strong interactions has improved in the last years, and we today have a candidate theory, QCD, phenomenological models still are the only way in which we can describe the ordinary multiparticle production in hadronic collisions.

In this thesis we study two such models: The hydrodynamic model introduced by Landau in 1953 [1], and a model developed here at Lund for baryon fragmentation.

The dominant feature of hadronic collisions that most of the secondary particles go into rather narrow cones in the forward and backward directions in the cms, was discovered rather early, in the 40's and the 50's.

Since the number of particles produced is rather large, statistical treatment of this multiparticle production was soon proposed, and in 1950 Fermi [2] constructed the first model of this kind. The main idea of this model is that a large amount of energy is deposited in a very small volume in a very short time. The volume is taken to be of the order of a Lorentz-contracted nucleon. Thermal equilibrium is assumed to be established at a high temperature, and particles emitted in analogy with blackbody radiation. This gives a mean number of particles $\langle n \rangle \sim s^{1/4}$ (where $\sqrt{s}$ is the energy in the cms), which agrees with data.

Landau argued in 1953 that this initially hot "prematter" should be treated as a relativistic fluid, and that it would expand predominantly in the longitudinal direction due to the large pressure in this direction. The expansion process, treated by relativistic hydrodynamics, then continues until the energy density equals that of a free hadron and the fluid breaks up into secondary particles.

The resultant distribution of the particles is a Gaussian distribution in rapidity, a velocity measure with simple transformation properties under Lorentz transformations.

During the 1950's the model was further developed by Russian and Japanese theorists. The only relevant data available at that time however came from exposures of emulsions to cosmic rays (i.e. protons colliding with nuclei) and this afforded only a rather preliminary comparison with data.

In the beginning of the 70's new high energy accelerators were put into operation: the ISR at CERN and the fixed target machine at Fermilab. Comparison with data could now be made more precisely. In the meantime the ideas of how multiparticle production is generated had changed to a large extent. Now so-called dual and multiperipheral models predicted a flat rapidity distribution in the central region. And the only effect of a larger incoming energy should be that this region extended further out in rapidity, giving a characteristic energy independent plateau in the distribution. But as experimental evidence grew and no such plateau with energy independent height was seen, the Landau hydrodynamic model (LHD) was reconsidered. In the LHD the experimentally observed rise of the central region with energy is a natural effect, and it was found that the agreement between data and the LHD was good. Since there are only approximate solutions to the equations of the model though, the fits are not wholly conclusive. The velocity of sound in the fluid is also usually taken as a free parameter to be determined by experiment and this also gives some extra freedom to the fits.

However, the success of the quark model explaining not only systematics, but also what happens in deep inelastic scattering and e^+e^- -annihilations, has focused the interest of particle physicists on explanations in terms of quarks and gluons, rather than on collective descriptions of statistical nature. Therefore interest

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The Effect of Temperature Dependence in the Hydrodynamic Model for Multiparticle Production.

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Abstract

The hydrodynamic model for particle production in high energy interactions is reconsidered for the case when the velocity of sound is allowed to vary with the temperature T of the fluid. The general equations of motion for the expansion phase are derived and a class of analytically solvable situations is considered in detail. This class includes systems close to free relativistic gas systems (corresponding to a temperature independent velocity of sound).

The general result is that the temperature dependence tends to flatten the rapidity distribution of the fluid elements, i.e. of the secondary particles. However, for the solutions close to free relativistic gas systems there are very small differences to the T -independent case. We conclude that the ordinary hydrodynamic model is stable with respect to the assumptions on the temperature dependence of the velocity of sound



1. Introduction.

The appearance of big accelerators at CERN and FNAL made it possible to study directly hadron-hadron interactions at high energies. Various models have been used to describe the resulting experimental data, among them the hydrodynamic model proposed by Landau in 1953 (LHD for short). [1] It is found that the agreement between the LHD and these data is fairly good at present energies. [2]

The physical picture that motivated Landau to propose his model was, however, rather different from our current view of hadronic collisions. Therefore it is of some interest to reconsider the basic assumptions of the LHD, to see whether or not they are compatible with our present knowledge of hadronic processes. Some of this work has been done [3] and in the present paper we give another piece of the program. We start with a brief account of the basic assumptions of the LHD (for a more complete version see [1]):

1. Initially, when two hadrons collide, a large amount of energy is deposited in a small region of space in a very short time. The interaction is assumed to be sufficiently strong to completely stop the matter momentarily, and a thermodynamic equilibrium is established.
2. Because of the large compression in the longitudinal direction, the matter expands and cools. During this phase it is assumed that the matter behaves as an ideal relativistic fluid.

3. When the density of the matter again reaches the normal hadronic density, the fluid breaks up and real particles, hadrons, appear.

The assumptions on the initial state have been criticized by Teper and Moravcsik [4], who found that if the matter were to stop completely, most of it would have disappeared as bremsstrahlung before equilibrium is reached. This indicates that if hydrodynamic models are to be taken seriously, the initial state must correspond to matter still in motion. Such a generalization of the initial conditions of the LHD is examined in [3]. The physical picture motivating that investigation is e.g. the quark-gluon models of the van Hove-Pokorski-type [5], where the glue, i.e. the part described hydrodynamically can be thought to be "dragged along" by the quarks. The results of the investigation are that even with an initial state with fairly slow-moving glue, the Gaussian rapidity spectrum of the LHD is flattened considerably. The so-called scaling solutions of the hydrodynamic equations, which automatically give a flat rapidity spectrum, also require a non-trivial initial state of motion [6]. Therefore we conclude in [3] that the predictions of the hydrodynamic model are unstable with respect to variations of the assumptions on the initial state.

We now turn to the expansion properties of the fluid. They depend on the assumptions of the equation of state. These assumptions occur in the fluid equations in connection with the velocity of sound c_1 of the fluid defined by

$$c_1^2 = \frac{dp}{d\varepsilon} \quad (1)$$

where p is the pressure and ϵ the energy density in the rest frame of the fluid. In his original work, Landau assumed the mass of the constituents to be negligible compared to the temperatures reached during the collision, and so chose c_1 to be $1/\sqrt{3}$, the value for a free relativistic gas of massless particles. Later works have considered the more general case when c_1 is a constant between 0 and 1. Comparison with experimental data then seems to favour values of c_1 between 0.4 and 0.5 [2c].

In this paper we will study the case when c_1 is allowed to vary with temperature. This is motivated if one assumes that the "prematter" that constitutes the fluid is made up of "particles", whose effective mass is not negligible compared to the thermal energy of the fluid. Since at present accelerator energies the temperatures reached in the model are not that high ($\lesssim 1$ GeV) [7], a nonvanishing mass of the constituents could be important. The results of this investigation are that the temperature dependence makes the final particle distribution flatter. However, the solutions that well describe a free relativistic gas, are those that do not give a markedly different particle spectrum. Thus the LHD is fairly insensitive to reasonable temperature dependent variations of the dynamical properties of the fluid.

The paper is organized in the following way: In section 2 we exhibit the equations of motion for the general case with arbitrary temperature dependence of the velocity of sound. We show that these equations can be solved for a special class of non-trivial situations. The thermodynamic properties of

these solutions are examined in section 3. We find that they in most cases behave qualitatively like a free relativistic gas. Section 4 then exhibits the final particle distributions and finally in section 5 our results are discussed and we summarize the conclusions. In an appendix we give the explicit derivation of the equations of motion and solve them in the special case.

2. A solution of the equations of motion.

In this section we give a solution of the equations of motion for the fluid, when c_1 is allowed to vary with the temperature T in the fluid. All the relevant quantities of the fluid can be obtained from its energy-momentum tensor $T^{\mu\nu}$. For an ideal fluid $T^{\mu\nu}$ only depends on the temperature T and the velocity $\bar{v}$ of the fluid elements [8]. To determine $T^{\mu\nu}$ we have to solve the equations of motion

$$\partial_\nu T^{\mu\nu} = 0 \quad (2)$$

Remembering that the expansion is mostly in the longitudinal direction we restrict ourselves in the following to 1+1-dimensions. Instead of T and v it is useful to introduce the variables κ and η defined by

$$\begin{cases} \kappa = \int \frac{\tau d\tau'}{c_1} & ; \tau = \ln T \\ \eta = \arctan(v) \end{cases} \quad (3)$$

Thus κ is another measure of the temperature and η the rapidity of the fluid element. Details of the calculation are left for the appendix. Here we just give the results.

Interchanging dependent and independent variables in the equations (2) we get two solutions just as for $c_1 = \text{const}$ (cf. [9]):

1. The progressive waves (along the characteristics)

$$x = \Gamma^\pm(\eta) \cdot t + f^\pm(\eta) \quad (4)$$

where

$$\Gamma^\pm(\eta) = \frac{sh\eta \pm c_1 \cdot ch\eta}{ch\eta \pm c_1 \cdot sh\eta} \quad (5)$$

$f^\pm(\eta)$ are arbitrary functions to be specified by the initial state.

2. The non-trivial solution where

$$\begin{aligned} x &= e^{-\tau} \{ sh\eta \cdot \partial_\tau \Psi - ch\eta \cdot \partial_\eta \Psi \} \\ t &= e^{-\tau} \{ ch\eta \cdot \partial_\tau \Psi - sh\eta \cdot \partial_\eta \Psi \} \end{aligned} \quad (6)$$

Here Ψ is a potential satisfying the equation

$$\partial_{\eta\eta}^2 \Psi - \partial_{\alpha\alpha}^2 \Psi = \Phi(\alpha) \cdot \partial_\alpha \Psi \quad (7)$$

with

$$\Phi(\alpha) = \frac{1}{c_1} \left[1 - c_1^2 - \frac{dc_1}{d\alpha} \right] \quad (8)$$

In deriving this equation it is essential that c_1 does not depend on η (see appendix).

For $\Phi(\kappa) = \text{const} = \Phi$ we can solve eq.(7). Expressing it in terms of the characteristic variables

$$\begin{cases} \rho = -\eta - \eta \\ \lambda = -\eta + \eta \end{cases} \quad (9)$$

the solutions is ^[10] (with the boundary determined by $\rho=0$ and $\lambda=0$)

$$\psi(\rho, \lambda) = \hat{\psi}(\rho, \lambda) \cdot \exp\left\{\frac{\Phi}{4}(\rho + \lambda)\right\} \quad (10)$$

where

$$\hat{\psi}(\rho, \lambda) = \hat{\psi}(0, 0) \cdot I_0\left(\frac{\Phi}{2}\sqrt{\rho\lambda}\right) + \int_0^\rho I_0\left(\frac{\Phi}{2}\sqrt{\lambda(\rho-r)}\right) \cdot \hat{\psi}_r(r, 0) dr + \int_0^\lambda I_0\left(\frac{\Phi}{2}\sqrt{\rho(\lambda-l)}\right) \cdot \hat{\psi}_l(0, l) dl \quad (11)$$

Here I_0 is the modified Bessel function of zeroth order and $\hat{\psi}_\rho$ and $\hat{\psi}_\lambda$ mean derivatives along the respective characteristic. Thus the solution is determined by the functions $\hat{\psi}_\rho(\rho, 0)$ and $\hat{\psi}_\lambda(0, \lambda)$. These functions are given by the initial conditions, and can be expressed in terms of the distributions $f^\pm(\eta)$ in eq.(4) from a comparison on the boundary of the non-trivial motion as

$$\begin{aligned} \partial_\rho \psi(\rho, 0) &= \frac{1}{2} e^{\frac{\Phi}{4}\rho} f^+(\eta) \cdot [c h \eta + c_1 \cdot s h \eta] \\ \partial_\lambda \psi(0, \lambda) &= -\frac{1}{2} e^{\frac{\Phi}{4}\lambda} f^-(\eta) \cdot [c h \eta - c_1 \cdot s h \eta] \end{aligned} \quad (12)$$

Therefore, given the distributions $f^\pm$, we have for $\Phi = \text{const}$ a class of solutions for the temperature dependent case.

3. The properties of the class of solutions, $\Phi = \text{const.}$

We will in this section determine the thermodynamic properties for a fluid with temperature dependent velocity of sound in the case when the function Φ in eq.(8) is a constant. This means that κ and c_1 are related by

$$\frac{dc_1}{d\alpha} = 1 - c_1^2 - c_1 \Phi \quad (13)$$

We note that if c_1 is independent of κ (i.e. of T), we get

$$c_1 = -\frac{\Phi}{2} \pm \sqrt{\frac{\Phi^2}{4} + 1} \quad (14)$$

Introducing the new constant

$$\varphi = \frac{\Phi}{\sqrt{\Phi^2 + 4}} \quad (15)$$

then gives

$$c_1 = \sqrt{\frac{1-\varphi}{1+\varphi}} \equiv c_+ \quad (16)$$

(the other solution is negative and thus not physical). If we now assume c_1 to be an increasing function of T

$$\frac{dc_1}{d\alpha} \geq 0 \quad (17)$$

c_1 becomes limited to

$$0 < c_1 < c_+ \quad (18)$$

As will be seen, the introduction of a break-up temperature will restrict the lower limit further.

We now turn to the explicit thermodynamic relations. The solution of eq.(13) gives:

$$\alpha = \alpha_0 + \frac{\sqrt{1-\varphi^2}}{2} \cdot \ln \left[\frac{(\sqrt{\frac{1+\varphi}{1-\varphi}} + c_1)}{(\sqrt{\frac{1-\varphi}{1+\varphi}} - c_1)} \right] \quad (19)$$

where κ_0 is an integration constant.

If we instead formulate eq.(13) in terms of τ by using $d\kappa = d\tau/c_1$ and resolve it, we find the relation between the temperature T and c_1 :

$$\left(\frac{T_0}{T}\right)^2 = \left(\sqrt{\frac{1-\varphi}{1+\varphi}} - c_1\right)^{1-\varphi} \cdot \left(\sqrt{\frac{1+\varphi}{1-\varphi}} + c_1\right)^{1+\varphi} \quad (20)$$

Here T_0 is an integration constant which we choose as the break-up temperature, at which real particles occur. In fig. 1 we exhibit c_1 as a function of T/T_0 . We notice that for $\Phi \gtrsim 2$ the temperature dependence is completely negligible, and then slowly becomes more and more pronounced as Φ decreases. Also given is $c_1(T)$ for a free relativistic gas with constituent mass m [11] plotted for different limiting temperatures T_0 . Note that corresponding to $T=T_0$, we have for each Φ a certain value $c_1(T_0)$, which we will call c_s , the velocity of sound at the break-up. As T tends to infinity on the other hand, c_1 tends to the limiting value c_+ . Thus we have (for finite temperatures)

$$c_s \leq c_1 < c_+ \quad (21)$$

c_+ and c_s as functions of Φ are shown in fig. 2.

The entropy density s of the fluid can be determined by using eq.(1) expressed in terms of T and s as

$$c_1^2 = \frac{s}{T} \frac{dT}{ds} \quad (22)$$

Equation (13) thus reformulated then gives

$$\frac{s}{s_0} = \frac{c_1}{\left(\sqrt{\frac{1-\varphi}{1+\varphi}} - c_1\right)^{\frac{1+\varphi}{2}} \cdot \left(\sqrt{\frac{1+\varphi}{1-\varphi}} + c_1\right)^{\frac{1-\varphi}{2}}} \quad (23)$$

where s_0 is an integration constant, which is fixed by the break-up condition. Figure 3 shows s as a function of T , where our normalization is such that s_c is the entropy at break-up, i.e. $s_c \equiv s(T_0)$. Here the temperature dependence of c_1 shows up as departures for low temperatures from the $s \sim T^{1/c_+^2}$ - behaviour one has for the corresponding temperature independent case.

For completeness we also give the energy density ε in terms of s , T and κ . This is obtained by using the relation $d\varepsilon = Tds$, and we find:

$$\varepsilon = \varepsilon_0 + \frac{1}{\phi^2 + 4} \left\{ sT \left(1 + \frac{\phi}{2c_1} \right) + s_0 T_0 (\mathcal{H} - \mathcal{H}_0) \right\} \quad (24)$$

(ε_0 another integration constant). Thus the value of ϕ determines both the thermodynamic properties in the model and the expansion properties of the fluid.

4. The final distribution of secondaries.

The observable quantity of interest is the inclusive spectrum of secondary particles in rapidity space, dN/dy . For hydrodynamic models this is proportional to the distribution of entropy, dS/dy , at the break-up temperature $T_c \approx m_\pi^{[1]}$.

Along the isotherm $T=T_c$ we have [6]

$$dS = s_\mu d\sigma' = s_c (ch y \cdot dx - sh y \cdot dt) \quad (25)$$

Using the Khalatnikov relations, eq.(6), and the potential eq.(7), we get

$$\frac{dS}{dy} = s_c \cdot A(y) \quad (26)$$

where

$$A(y) = c_1^2 \cdot e^{-\tau} \cdot \partial_\tau U \Big|_{\tau=\tau_c} \quad (27)$$

with

$$U \equiv \Psi - \partial_\tau \Psi \quad (28)$$

We have here assumed that a particle will appear with the same rapidity as that of the fluid element that it stems from (i.e. $y=\eta$). Thus once we know the potential Ψ we can compute the distribution dS/dy . The only thing we have to do is to specify the initial functions $f^\pm$.

In the following we will restrict ourselves to the initial conditions of the Landau model, where the matter expands from a completely uniform distribution, and the $f^\pm$ are given by the longitudinal width of the distribution. Thus we may choose

$$\begin{cases} f^+ = 0 \\ f^- = d \end{cases} \quad (29)$$

where d is a length of the order of the width of a contracted hadron, and pursue to investigate the consequences of our temperature dependent velocity of sound solutions.

Before doing so, however, we must determine the normalization. The Landau model contains two temperatures of importance: the temperature T_0 at break-up but also the temperature T_M of the initial state. T_M is the maximum temperature reached in the system. In the following we need to specify κ at these

two temperatures. We do this in analogue with c_1 -temperature independent calculations e.g. [3], choosing

$$\begin{aligned} \mathcal{K}(T_M) &= 0 \\ \mathcal{K}(T_0) &= -\eta_0 \end{aligned} \quad (30)$$

The parameter η_0 is the rapidity of the fastest fluid elements assumed to end up as secondary particles, and thus determines the width of the dS/dy distribution.

We now proceed by inserting eqs.(29) in eqs.(12) and integrating, remembering that

$$\tau'(\lambda) = \frac{d\tau}{d\mathcal{K}} \cdot \frac{d\mathcal{K}}{d\lambda} = -\frac{1}{2} c_1(\lambda) \quad \text{for } \mathcal{S}=0 \quad (31)$$

and fixing the integration constant by choosing $\Psi(0,0) = 0$.

In this way we get

$$\begin{aligned} \Psi(\mathcal{S}, 0) &= 0 \\ \Psi(0, \lambda) &= -\frac{d}{2} \left\{ e^{\tau + \frac{\lambda}{2}} - e^{\tau - \frac{\lambda}{2}} \right\} \end{aligned} \quad (32)$$

Equation (11) then gives the potential

$$\hat{\Psi}(\mathcal{S}, \lambda) = \frac{d}{2} \int_0^{\lambda} I_0\left(\frac{\Phi}{2} \sqrt{\mathcal{S}(\lambda-l)}\right) \left\{ p(l) e^{P(l)} - q(l) e^{Q(l)} \right\} dl \quad (33)$$

where

$$\begin{aligned} P(\lambda) &= \tau(\lambda) + \frac{\lambda}{2} - \frac{\Phi}{4} \lambda \\ Q(\lambda) &= \tau(\lambda) - \frac{\lambda}{2} - \frac{\Phi}{4} \lambda \end{aligned} \quad (34)$$

and

$$\begin{aligned} p(\lambda) &= -P'(\lambda) = \frac{1}{2} c_1(\lambda) - \frac{1}{2} + \frac{\Phi}{4} \\ q(\lambda) &= -Q'(\lambda) = \frac{1}{2} c_1(\lambda) + \frac{1}{2} + \frac{\Phi}{4} \end{aligned} \quad (35)$$

Working the algebra out and using eq.(8) and the normalization,

eq. (30), we obtain:

$$\begin{aligned}
 A(\eta) = & \frac{1}{T_0} e^{\frac{\phi}{2} \eta_0} \frac{d}{2} \left\{ \left(\frac{1}{c_s} - c_1(T_M) - \frac{\phi}{2} \right) \cdot T_M \cdot I_0 \left(\frac{\phi}{2} \sqrt{\eta_0^2 - \eta^2} \right) + \right. \\
 & + T_M \frac{\phi}{2} \frac{\eta_0}{\sqrt{\eta_0^2 - \eta^2}} I_1 \left(\frac{\phi}{2} \sqrt{\eta_0^2 - \eta^2} \right) - \\
 & - \left(1 + \frac{\phi}{2 c_s} \right) \int_0^{\eta_0 + \eta} I_0 \left(\frac{\phi}{2} \sqrt{(\eta_0 - \eta)(\eta_0 + \eta - l)} \right) \left[p(l) e^{P(l)} - q(l) e^{Q(l)} \right] dl - \\
 & \left. - \frac{1}{c_s} \int_0^{\eta_0 + \eta} I_0 \left(\frac{\phi}{2} \sqrt{(\eta_0 - \eta)(\eta_0 + \eta - l)} \right) \left[(q(l) + \frac{\phi^2}{8}) e^{Q(l)} + (p(l) - \frac{\phi^2}{8}) e^{P(l)} \right] dl \right\} \quad (36)
 \end{aligned}$$

For the case $c_1 = \text{const} = c_0$ this formula reduces to that of the Landau model, as it should:

$$A(\eta) \Big|_{c_1 = \text{const}} = \frac{T_M}{T_0} \frac{d}{2} e^{\frac{\phi}{2} \eta_0} \frac{\phi}{2} \left\{ I_0 \left(\frac{\phi}{2} \sqrt{\eta_0^2 - \eta^2} \right) + \frac{\eta_0}{\sqrt{\eta_0^2 - \eta^2}} I_1 \left(\frac{\phi}{2} \sqrt{\eta_0^2 - \eta^2} \right) \right\} \quad (37)$$

In fig. 4 we exhibit $A(\eta)$ for different values of ϕ . When c_1 is temperature independent, $\phi = (1 - c_+^2)/c_+$. So to see the effect of c_1 being temperature dependent, we also show the ordinary LHD-spectra for the corresponding values of $c_1 = c_+$. Notice that for $\phi \gtrsim 2$, i.e. c_1 practically temperature independent, as seen in fig. 1, we have the normal Gaussian shape of the distribution. As ϕ decreases, the distribution flattens until for $\phi \lesssim 1$ we get a rising spectrum. Since we do not expect this to be the case, $\phi = 1$ thus stands as a limit for what is physically reasonable.

5. Discussion and concluding remarks.

In this paper we have studied the influence of a temperature dependent velocity of sound in the LHD. We have derived the general equations and solved them in a special case, parametrized by Φ . As can be seen in fig. 4 the final particle distributions for these T-dependent solutions are flatter than the corresponding T-independent ones.

The general shape of $c_1(T)$ for the solutions we have found is very much like that of a free relativistic gas whose constituents have mass m [11]. From fig. 1 we see that the best fit is around $T_0 = 0.4 m$ and $\Phi = 1.2$. If $T_0 \approx m_\pi$ this gives $m \approx 350$ MeV which seems too big to be realistic. However, the solutions are certainly much better fitted by an inhomogenous gas. For instance if we assume that the gas has a mass spectrum proportional to m^a (a being a constant) for m greater than some limiting value m_0 , $c_1(T)$ will tend to $\frac{1}{\sqrt{a+4}}$ when $T \rightarrow \infty$ [2a]. Comparing our solutions $c_1(T)$ to a Bose-Einstein gas with this kind of mass spectrum ^{*/} then gives good agreement (less than 1% per point) for values of $a \gtrsim -0.5$ and $m_0 \lesssim (a + 1.5)m_\pi$. However, we note that (see fig. 4), if $\Phi \gtrsim 1.3$ (corresponding to $a \gtrsim -0.5$) the effect of the temperature dependence hardly changes the final particle distribution at all.

In conclusion we see that assumptions about the constitution of the fluid, i.e. about the equation of state, eq.(1),

^{*/} A Fermi-Dirac gas gives almost the same result.

naturally are important for the outcome of the model. For the LHD this investigation has however shown that the final particle distribution is surprisingly insensitive to the temperature dependence of c_1 for those solutions that are close to free relativistic gases with not too large an internal mass scale. If we, on the other hand, accept the lower mass limit m_0 in the gas to be around 350 MeV, temperature dependence becomes important. It is noteworthy that this value is not too far from that usually expected for the scale parameter Λ in QCD, the tentative theory of strong interactions.

Appendix

We give in this appendix the derivation of eq. (7) and the solution, eq.(10), presented in section 2.

The equations of motion (2) for an ideal fluid in 1+1-dimensions are ^[9] in terms of the characteristic variables eq.(9):

$$\begin{aligned} (1+v \cdot c_1) \frac{\partial \rho}{\partial t} + (v+c_1) \frac{\partial \rho}{\partial x} &= 0 \\ (1-v \cdot c_1) \frac{\partial \lambda}{\partial t} + (v-c_1) \frac{\partial \lambda}{\partial x} &= 0 \end{aligned} \tag{A1}$$

where v is the velocity of the fluid element. These equations, where ρ and λ are expressed as functions of x and t , are very hard to solve. In order to make any progress we interchange dependent and independent variables. I.e., we express x and t as functions of ρ and λ .

A necessary and sufficient condition for such a transformation to be valid is

$$D \equiv \begin{vmatrix} \partial_t \rho & \partial_x \rho \\ \partial_t \lambda & \partial_x \lambda \end{vmatrix} \neq 0 \quad (A2)$$

If on the other hand $D=0$, we get a special solution, the characteristic one giving the progressive waves. The reason is that $D=0$ means that we are on the characteristics. In our case these are expressed as $\rho = \text{const}$ or $\lambda = \text{const}$. Solving eqs.(A1) for $\rho=0$ and $\lambda=0$ respectively, exactly as it is done in [9] for $c_1=\text{const}$, then gives eqs.(4) and (5).

On the other hand, if $D \neq 0$ we get

$$\begin{aligned} \partial_\lambda t(c_1 \cdot ch\eta + sh\eta) - \partial_\lambda x(ch\eta + c_1 \cdot sh\eta) &= 0 \\ \partial_\rho t(c_1 \cdot ch\eta - sh\eta) + \partial_\rho x(ch\eta - c_1 \cdot sh\eta) &= 0 \end{aligned} \quad (A3)$$

Let us now introduce the expressions:

$$\begin{aligned} K &= \frac{1}{2} e^\tau \{ t(sh\eta - c_1 \cdot ch\eta) + x(c_1 \cdot sh\eta - ch\eta) \} \\ L &= \frac{1}{2} e^\tau \{ -t(sh\eta + c_1 \cdot ch\eta) + x(c_1 \cdot sh\eta + ch\eta) \} \end{aligned} \quad (A4)$$

As a consequence of eqs.(A3) we then get

$$\begin{aligned} \partial_\rho K &= \frac{1}{4} e^\tau \left[-1 + c_1^2 - 2 \frac{\partial c_1}{\partial \rho} \right] (t \cdot ch\eta - x \cdot sh\eta) \\ \partial_\lambda L &= \frac{1}{4} e^\tau \left[-1 + c_1^2 - 2 \frac{\partial c_1}{\partial \lambda} \right] (t \cdot ch\eta - x \cdot sh\eta) \end{aligned} \quad (A5)$$

And so we see that if

$$\frac{\partial c_1}{\partial s} - \frac{\partial c_1}{\partial \lambda} = \frac{\partial c_1}{\partial \eta} = 0 \quad (A6)$$

we have

$$\partial_\lambda L = \partial_s K \quad (A7)$$

This means that it is possible to introduce a potential

$\Psi \equiv \Psi(\rho, \lambda) \equiv \Psi(\eta, \tau)$, such that

$$\begin{cases} L = \partial_s \Psi \\ K = \partial_\lambda \Psi \end{cases} \quad (A8)$$

If we now from eq.(A4) solve x and t in terms of K and L , and then use eq.(A8) to express x and t in terms of the potential, we get the ordinary Khalatnikov relations, eqs.(6).

Since Ψ is a function of τ and η , eqs.(6) now express x and t as functions of τ and η . Thus, if we know Ψ we can compute all the relevant facts about the non-trivial expansion of the fluid.

In order to determine Ψ , we form the following equation

$$\partial_s K + \partial_\lambda L = 2 \frac{\partial^2 \Psi}{\partial s \partial \lambda} = \frac{1}{4} e^\tau \left[2(1-c_1^2) - 2c_1 \frac{dc_1}{d\tau} \right] (x \cdot sh \eta - t \cdot ch \eta) \quad (A9)$$

But

$$L + K = \partial_s \Psi + \partial_\lambda \Psi = -\partial_x \Psi = e^\tau c_1 (x \cdot sh \eta - t \cdot ch \eta) \quad (A10)$$

then implies

$$\frac{\partial^2 \psi}{\partial s \partial \lambda} = \frac{1}{c_1} \left[1 - c_1^2 - c_1 \frac{dc_1}{d\tau} \right] (\partial_s \psi + \partial_\lambda \psi) \quad (A11)$$

or reformulated in terms of κ and η we have eq.(7).

To solve eq.(A11), we introduce

$$\psi = \hat{\psi} \cdot \exp \left[-\frac{1}{2} \int d\alpha' \Phi(\alpha') \right] \quad (A12)$$

Reformulating eq.(7) in terms of $\hat{\psi}$ then gives

$$-\partial_{\eta\eta}^2 \hat{\psi} + \partial_{\alpha\alpha}^2 \hat{\psi} = \frac{1}{4} \Lambda \cdot \hat{\psi} \quad (A13)$$

where

$$\frac{1}{4} \Lambda = \frac{1}{4} \Phi^2(\alpha) + \frac{1}{2} \Phi'(\alpha) \quad (A14)$$

This is a 2-dimensional Klein-Gordon equation, where $\Lambda/4$ plays the role of a mass term. We notice that if $\Phi(\kappa)$ is a constant, Λ becomes a constant equal to Φ^2 , and eq.(A13) is solvable. The Riemann function of eq.(A11) is then [10]

$$I_0 \left(\frac{\Phi}{2} \sqrt{(\lambda-l)(s-r)} \right) \quad (A15)$$

and integrating along $r=0$ and $l=0$ respectively then gives eq.(11). Thus we have arrived at the solution presented in section 2.

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Figure captions.

- Fig. 1 The velocity $c_1(T)$ for different values of the parameter Φ (full lines). Also shown is $c_1(T)$ for a free relativistic gas with constituent mass m [11] (broken lines).
- Fig. 2 c_s and c_+ as functions of Φ .
- Fig. 3 The entropy density s as a function of temperature T for different values of Φ .
- Fig. 4 Final particle distributions $A(y) \sim \frac{dN}{dy}$ for
a) LHD with $c_1 = \text{const} = c_+$,
b) LHD when c_1 is T-dependent and Φ varies.

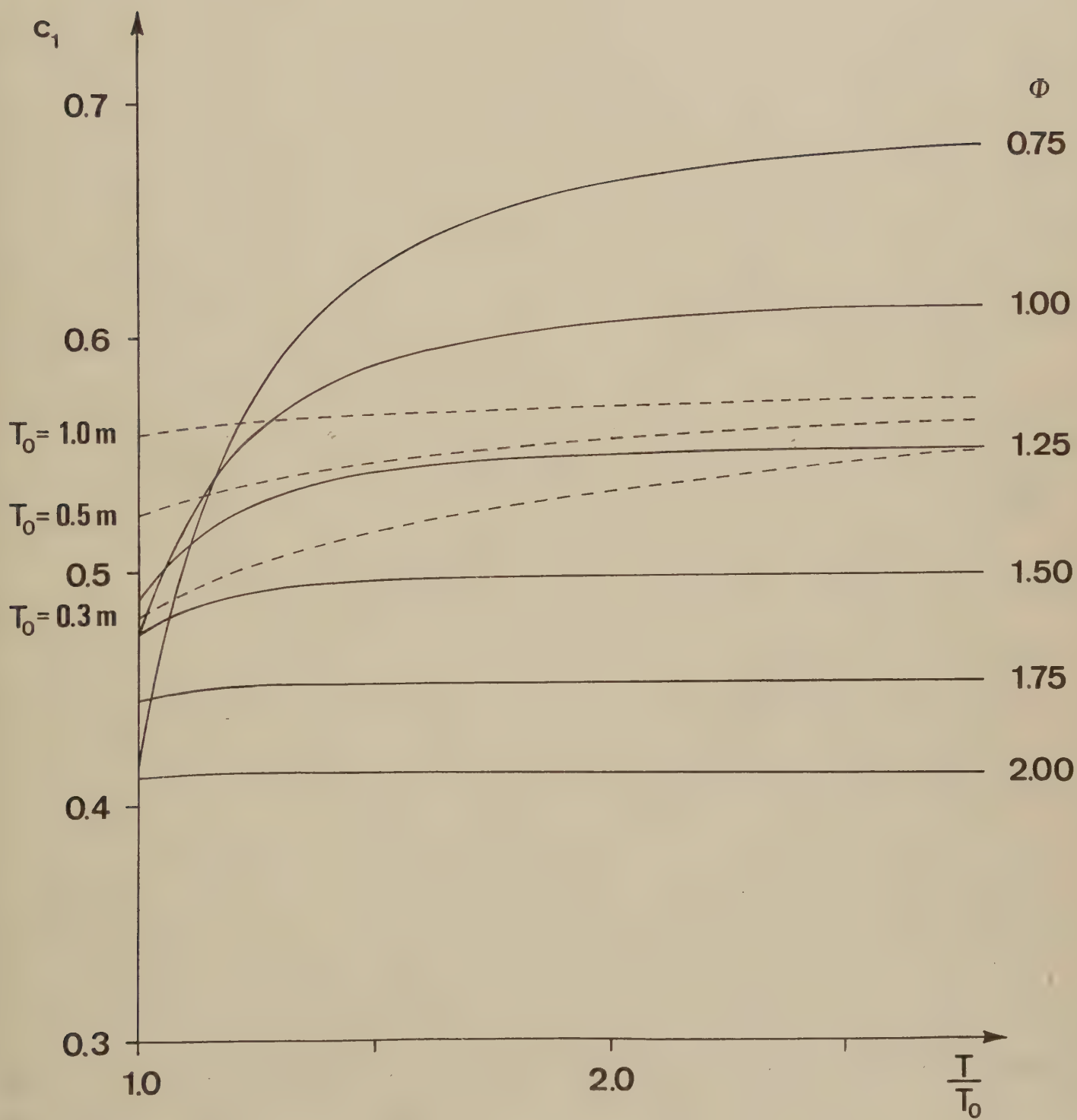


Fig. 1

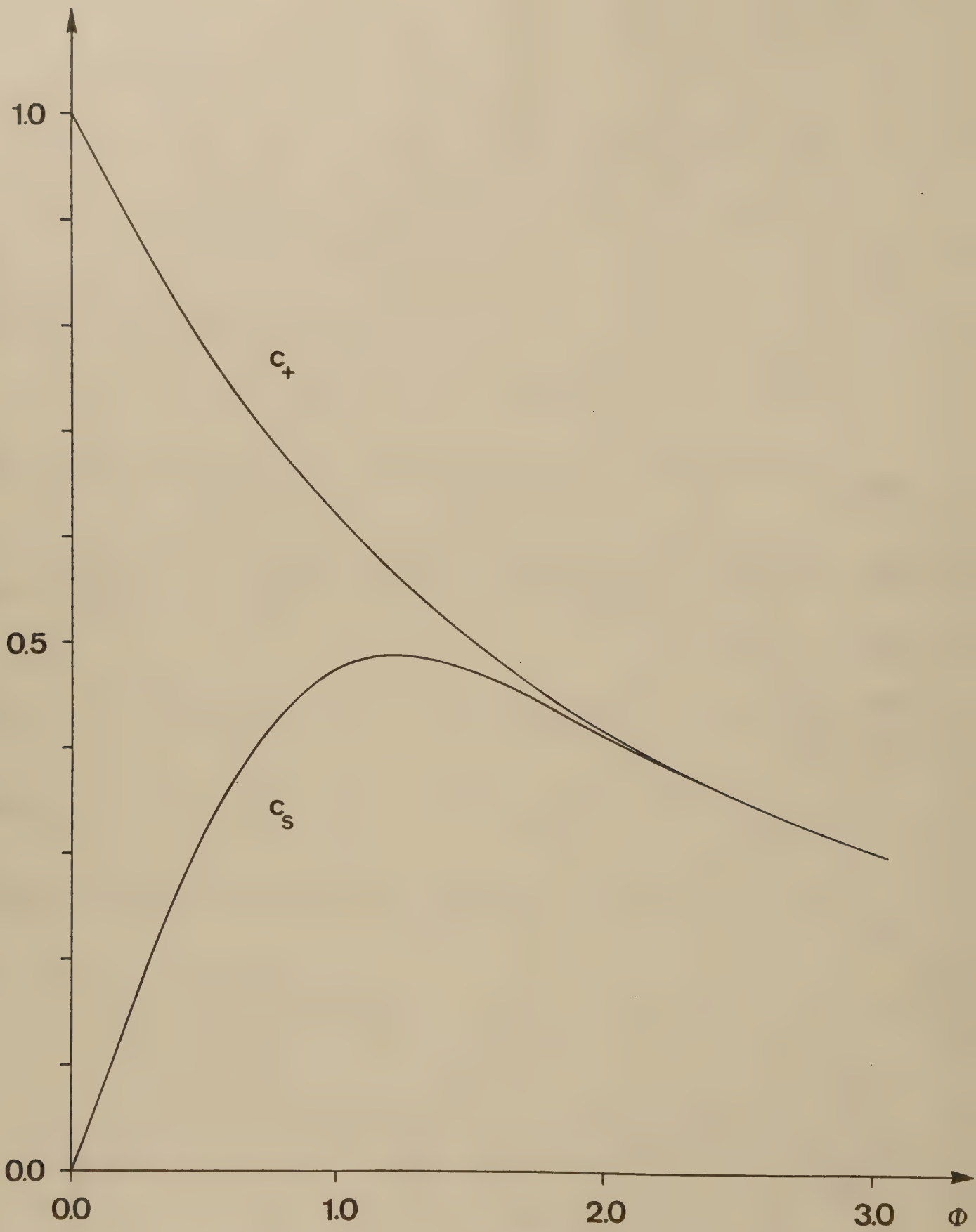


Fig. 2

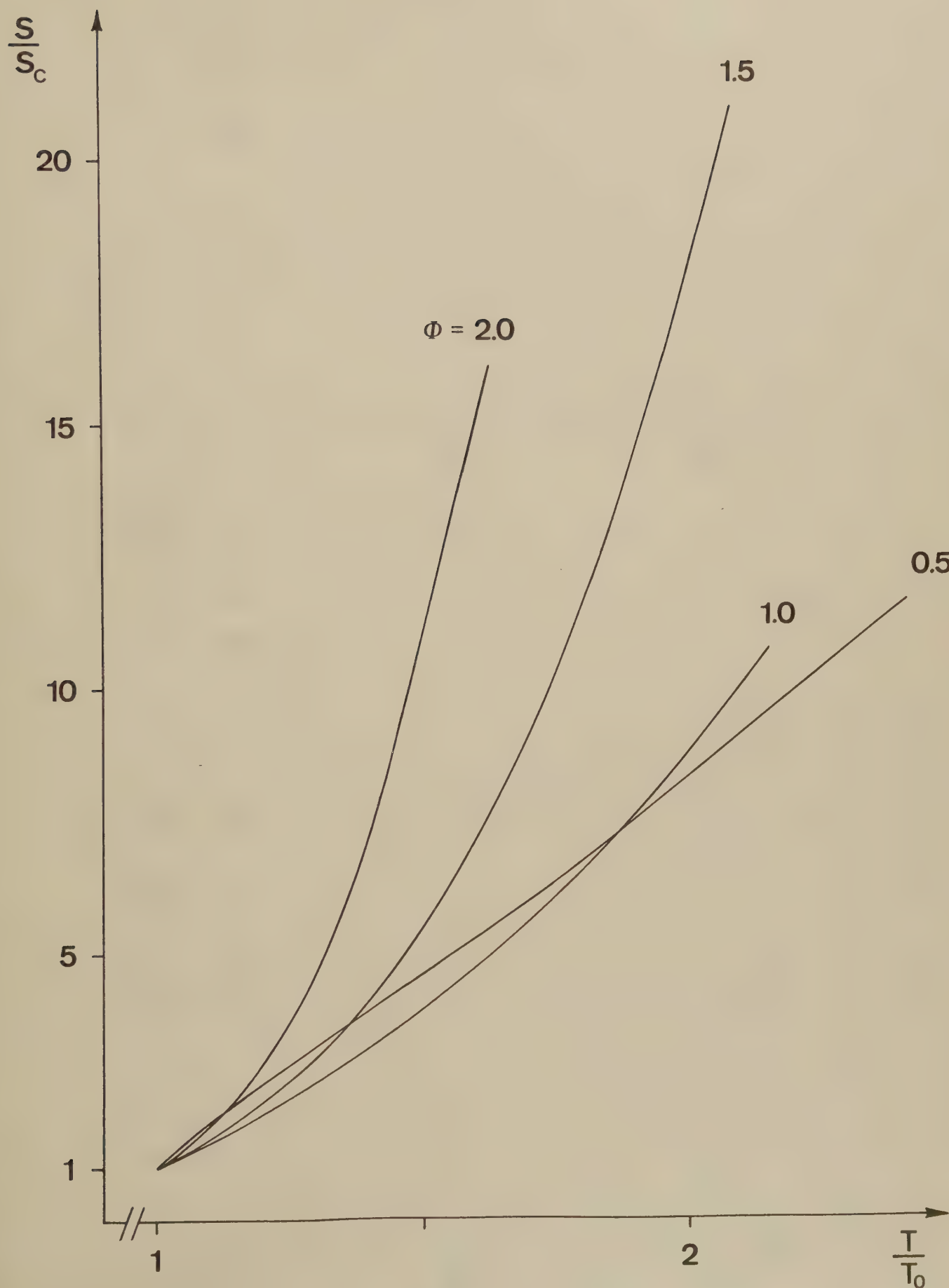


Fig. 3

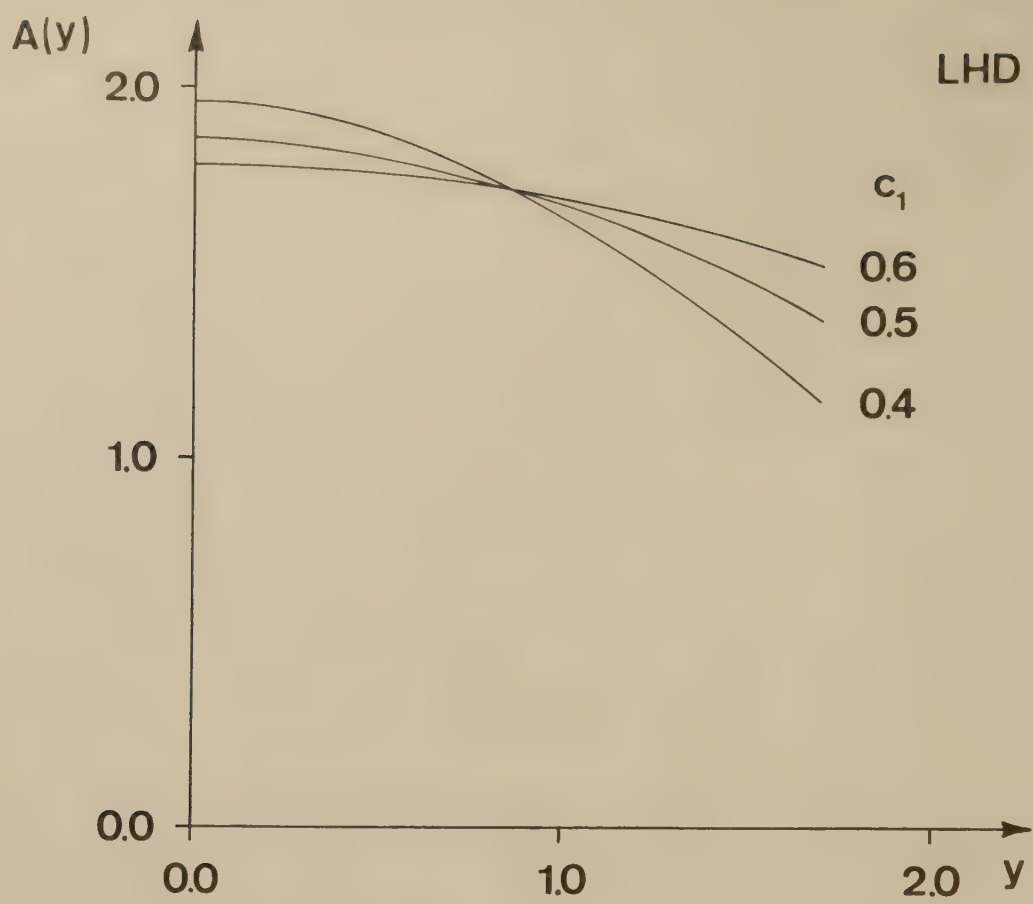


Fig. 4a

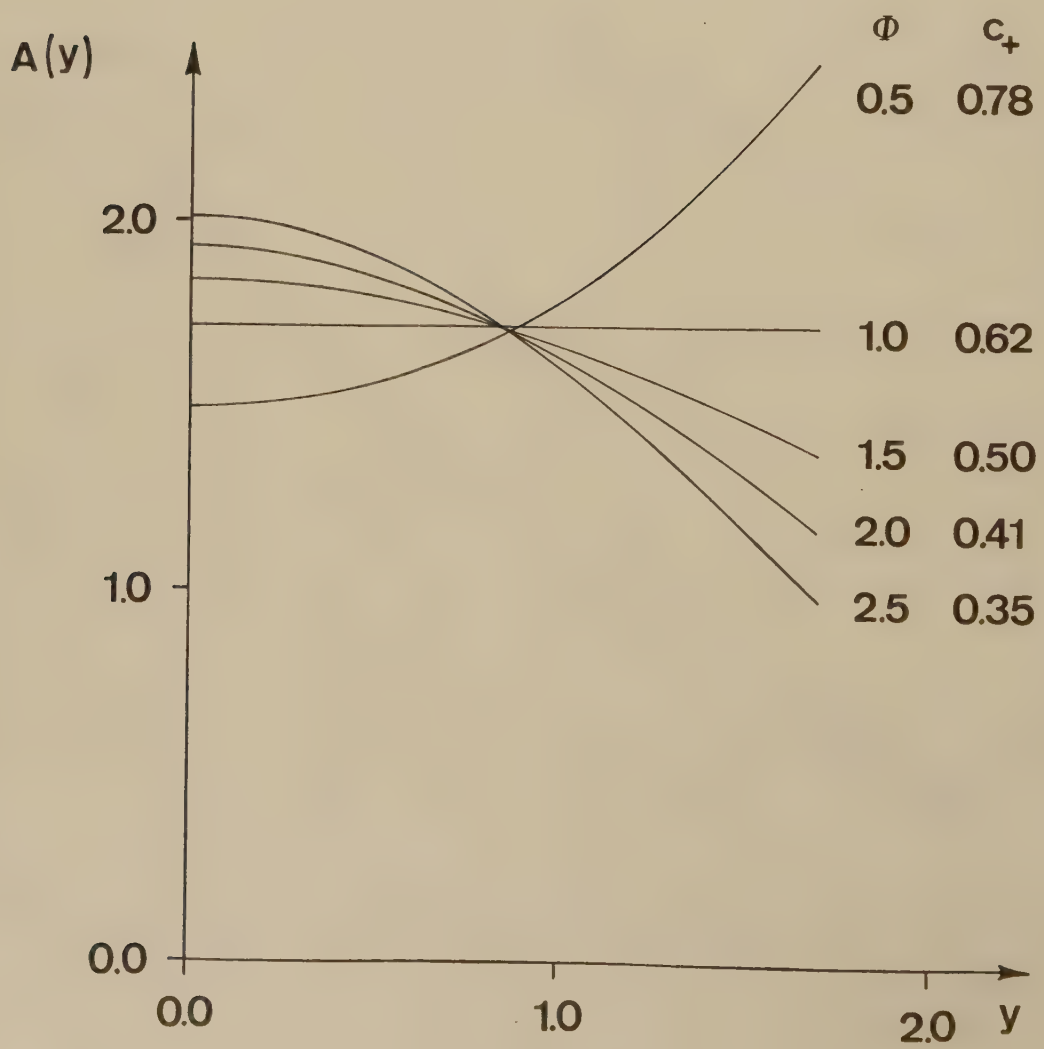


Fig. 4b

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A MODEL FOR THE REACTION MECHANISM AND THE BARYON FRAGMENTATION
DISTRIBUTIONS IN LOW $p_{\perp}$ HADRONIC INTERACTIONS

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Abstract: We present a reaction mechanism, based on a semiclassical treatment of color dynamics together with the corresponding results for the final state in a low $p_{\perp}$ hadronic interaction. The ensuing model is of fragmentation character (like the one presented on phenomenological grounds in refs [1,2]) and the final state hadrons stem from the breakup of an essentially one-dimensional color force field in a way similar to the models for soft hadronisation of colored objects that we have developed for use in leptonproduction and e^+e^- -annihilation. We present a set of predictions from the model for baryon fragmentation distributions in low $p_{\perp}$ reactions which is in good agreement with the present experimental data and also contains some nontrivial correlation properties that can be tested in the next generation of experiments.

1. Introduction

In this paper we present a complete model for the particle distributions in the fragmentation regions of low $p_{\perp}$ hadroproduction.

The intention is to obtain a model of "the parton fragmentation type" (like the one we have suggested on a purely phenomenological basis in refs. [1,2]), which on the one hand contains a simple and self-consistent dynamical scheme and on the other hand agrees with the published experimental data.

This paper is mainly concerned with a discussion of the baryonic fragmentation regions but the general formalism just as well applies to mesonic ones.

In section 2 we present a reaction mechanism for an inelastic low $p_{\perp}$ event, which is based on a semi-classical approach to color dynamics. The reaction mechanism will be further elaborated in a forthcoming work [3], in particular it will be used to compute cross sections, overlap functions etc. This reaction mechanism will on a long time development scale lead to a final state of an essentially one-dimensional character like a string with no excited transverse degrees of freedom (a "stretched-out color bag").

Both for an impinging meson and for a baryon, one of the valence constituents will be at the end of the force-field. However, for the baryon case, this leading quark will have another valence quark (a partner in a leading diquark) somewhere behind it along the force-field. Thus the diquark will in our work correspond to an object with some internal structure. In particular there is a certain probability, which we treat in a phenomenological way, that it breaks up into two quarks which end up in different final state hadrons.

As regards energy and momentum the resulting final state is similar to the colored force field which is stretched-out by an ordinary quark (antiquark) jet in leptonproduction or in e^+e^- annihilation. Therefore we expect strong similarities between the two situations in connection with the energy-momentum partition among the final state particles, although the dynamical processes are very different. The stretched-out color force-field will break up by the production of quark-antiquark ($q\bar{q}$) pairs along the field just as in the model for the soft hadronisation we have developed based on a semiclassical approach to string dynamics in refs. [4,5]. This model is causal and relativistically invariant and of a stochastic character, which readily lends itself to an implementation in terms of a Monte Carlo generation program [6].

Assuming the decay of a stretched-out color force field to be independent of the way it is created we will, with a few modifications due to the influence of the valence flavors which are not at the end-point of the force-field, use the same scheme. We present a few details of this scheme in sections 3 and 4. For baryon production in the baryonic fragmentation regions we use SU(6) symmetry, both to characterise the original baryon diquarks and also to determine the final state baryon quantum numbers. We will only take into account the production of baryons in the ground state 56-multiplet and similarly for the mesons only the ground state pseudoscalar and vector mesons.

We note that the ratio of vector mesons to pseudoscalar mesons produced in the field seemsto be around 1:1 [5] . One may

therefore expect that there is a corresponding suppression of spin $3/2$ states for the baryons. A phenomenological parameter for the production ratio of decuplet to octet baryons can easily be introduced in the model. There are indications (from Σ - and Σ^* -production rates) that this ratio is of the order 1 instead of 2, expected for equal population of all available spin states. The precise value of this parameter has only a small influence on the number of protons, neutrons and mesons.

We note in passing that the model contains a linear force field of the kind we assumed for the model of inclusive Λ -polarisation presented in ref. [7]. The main point is a correlation between the transverse momenta and the spin of $q\bar{q}$ -pairs generated along a force field with no excited transverse degrees of freedom. A similar mechanism also implies a nonvanishing polarisation of inclusively produced Ξ particles.

In section 5 we present a large amount of predictions for the single particle distributions as well as correlation distributions. We compare the predictions to the available experimental data wherever possible and find that not only the main features but also many small details are reproduced by the model predictions. An interesting measure of the total event shapes is introduced by Ochs and Stodolsky in ref. [8]. The model reproduces these jet profiles in good agreement with the experimental data.

We will in this paper only consider the longitudinal momentum distributions, leaving transverse momentum properties as well as polarisation properties for a future publication.

We end with a few concluding remarks in section 6.

2. The interaction mechanism

We picture the hadrons as extended objects, bags, containing color which is confined by the pressure from the surrounding vacuum. Due to the properties of asymptotic freedom the color of e.g. a quark is not a well localised quantity (see fig. 1), but rather smeared out over a distance of $\approx 1/\Lambda \approx 0.4$ fm like a colored "blob". Thus a meson bag contains two such blobs, one colored and one anti-colored, while a baryon bag correspondingly contains three color blobs.

In the ground state of a hadron the blobs will be partly overlapping (cf. fig. 1) and we expect that on the space-time scale relevant for low $p_{\perp}$ reactions this overlap region will contain a color force field but no color charges. This is in contrast to the situation in a leptonproduction event where a much smaller space time size is probed. The non-overlapping regions will, however, contain color (anticolor) charges and the hadron bag therefore looks as an object containing colored (anti-colored) "handles" surrounding a "white" centre.

We note that the vacuum pressure will be the same from all directions and therefore an isolated hadron bag can move freely without energy and momentum losses like a bubble in a non-viscous fluid. Although we do not need in this note a precise structure for the color fields we note that the resulting picture is compatible with a parton picture where the valence constituents are surrounded by wee partons.

When two objects of this kind collide there may be some overlap between the color handles from each bag. We then expect that with a certain probability there will be a connection between the blobs, i.e. the vacuum will not close again when the hadrons go apart. Instead the remnants will be connected via a color tunnel as in fig. 2 and the whole object becomes one stretched-out bubble in the vacuum.

We note that this implies that there is no longer any vacuum pressure behind the connected blobs. Therefore the result is that the pressure in front acts on the blobs just like an attractive force between them. We expect that the force is very similar to the one encountered in the color fields used in the soft hadronisation models of colored objects in leptonproduction [4,5] or in the potential models for charmonium spectroscopy [26] i.e. corresponding to an energy loss per unit length equal to the string constant $\kappa \sim 1 \text{ GeV/fm}$.

In a high energy collision the internal motion between the color blobs of each hadron is strongly time-dilated in e.g. the cms frame of the collision. Therefore the remaining valence blob(s) i.e. those without a color tunnel behind, will continue to move essentially unaffected until all the energy of the interacting blob is used up to produce the force field behind it. When this happens this blob is stopped and the vacuum pressure starts to work on the next one. The second blob will now be retarded with the same force and so on, until all the original hadron energy-momentum is used up.

In this way the interaction will on a longtime scale cause the (step wise) stretching out of an essentially one-dimensional color field in such a way that when all the energy momentum is used up, one of the valence flavors will be at each end point of the field. We note, in particular, that the result is the same whatever is the original partitioning of the hadron energy momentum among the partons.

In general, however, long before the energy momentum is used up, the color tunnel will break up by the production of new $(q\bar{q})$ pairs along the field as discussed in the next section.

3. The soft hadronisation mechanism

We will treat the stretched-out color tunnel resulting from the reaction mechanism of section 2 as a massless relativistic string without excited transverse degrees of freedom as we have done for the color fields in the semiclassical approach of refs. [4,5]. For completeness we will in this section briefly describe the resulting model for particle production. In the next section we introduce the modifications necessary to take into account the specific properties of the final state in a low $p_{\perp}$ reaction such as the extra valence flavor(s).

Consider the situation when a massless $(q\bar{q})$ pair, (subsequently called $q_0\bar{q}_0$) with large energy W , is produced in a single space time point and starts to move apart in opposite directions like in an e^+e^- -annihilation event. The particles will then move along the light cones stretching out a string-like force field

between them, thereby losing energy-momentum at the rate of $\kappa \approx 1 \text{ GeV/fm}$. Along the force-field new $(q\bar{q})$ pairs will be produced and afterwards be dragged apart by the force field, the q -particle in one direction (towards the anticolored end) and the $\bar{q}$ -particle in the opposite direction. In that way the force field will break up into small string pieces, each one containing one q - and one $\bar{q}$ -particle at the end points (cf figs 3,4).

In this model the mesons can be ordered according to flavor ("rank") because the produced $(q\bar{q})$ pairs have total flavor zero. Due to the stochastic nature of the production process flavor-ordering will, however, only in the mean reflect itself in rapidity ordering. We note, however, that ordering according to production time is a Lorentz frame dependent notion. In this causal and relativistically invariant scheme the slowest particles will be the first to get disentangled from the jet in each frame (an inside-out cascade).

Therefore in a frame moving very fast with respect to the cms along the q_0 -particle direction the first rank meson, containing the q_0 flavor, will also be produced first in time. The q_0 -particle will actually in this frame only have a small initial momentum and therefore it will be rapidly stopped and turned around by the force field to get accelerated after the $\bar{q}_0$ -particle (fig. 5). When the pair $(q_1\bar{q}_1)$ is produced, $(q_0\bar{q}_1)$ will form the first rank meson with a fraction z of the original q_0 jet energy momentum, related to the distance marked in fig. 5. From the results that the density of states for the break up

system, composed of the meson $(q_0 \bar{q}_1)$ and the remainder jet $(q_1 \bar{q}_0)$ with the large mass $M^2 \approx W^2(1 - z)$, is $\frac{dn}{dM^2} = \text{constant}$ in a linear force field and the assumption that all available final states have the same production rate, we obtain the result that the probability to obtain the meson $(q_0 \bar{q}_1)$ at z is constant, independent of z . Repeating the process we obtain a model of the iterative cascade type

When the model is extended to the production of heavy mass $(q\bar{q})$ pairs as well as pairs with transverse mass, $\mu_\perp$, in a constant force-field with no excited transverse degrees of freedom we obtain a production rate factor for each vertex

$$e^{-\left[\frac{\mu_\perp^2 \pi}{\kappa}\right]} |g|^2 |h|^2 d\mathbf{p}_\perp^2 \quad (1)$$

This factor implies in particular that the production of $(s\bar{s})$ pairs should be suppressed as compared to $(u\bar{u})$ pairs and $(d\bar{d})$ pairs with a factor of $1/3 - 1/4$ (i.e. the usually assumed SU(3) breaking parameter [9]) and that the transverse momentum is locally compensated in each $(q\bar{q})$ pair and has a Gaussian distribution with a mean $\sim (.30 - .35)$ GeV/c (which is also generally assumed [9]).

The factor $|g|^2 = |g|^2((\kappa\tau)^2/\mu_\perp^2)$ with τ the invariant distance from the original production point of the $(q_0 \bar{q}_0)$ pair to the production point of the pair with transverse mass $\mu_\perp$ (cf. fig. 6) has the following interpretation. Classically it is necessary in order to produce a state with the energy $2\mu_\perp$ from the field energy to have a smallest field length $\frac{2\mu_\perp}{\kappa}$. We have found that

this corresponds to a quantum mechanical suppression (described by $|g|^2$) unless the effective field length 2τ is much larger than $\frac{2\mu_{\perp}}{\kappa}$. The model independent features of $|g|^2(\alpha)$ are that it is small for small α , rapidly increases to 1 for $\alpha \gg 1$ and exhibits a smooth behavior in between. The above-mentioned properties imply a noticeable suppression of particles with $z \sim 1$ and a stronger correlation between transverse and longitudinal momenta than the one implied by pure energy momentum conservation. However, the precise features of $|g|^2$ do not reflect themselves in the final state particle distributions.

While the factor $|g|^2$ has a strong both theoretical and experimental support the second factor $|h|^2$ in eq.(1) is less strongly motivated. In most applications to quark and gluon jets, and also for those in this paper, the factor can be neglected and put equal to 1, but in some minor parts of phase-space and for some particular flavor quantum number combinations the effect of $|h|^2$ is noticeable. To understand the motivation for $|h|^2 = |h|^2((\kappa\tau')^2/\mu_{\perp}^2)$ consider fig. 7. Here one production vertex for a heavy (transverse) mass pair V' is very "late", compared to the adjacent vertices V_1 and V_2 , so that one meson is followed in rank by another one having much more energy. The occurrence of such a situation is rare, due to the competition from the many "earlier" possible vertex points in the stochastic process.

In a classical description the production of the pairs at the vertices marked V_1 and V_2 in fig. 7 implies that the force

field vanishes outside the light cone segments $V_1 O'$ and $V_2 O'$. Therefore the effective field length $2\tau'$ (with τ' the invariant distance between O' and V') is smaller than the one needed to produce the $(q\bar{q})$ pair at rest $\frac{2\mu_{\perp}}{\kappa}$ and one may expect a suppression factor similar to the one obtained in connection with fig. 6, i.e. $|g|^2$. We note, however, that in this case it is also necessary to understand the quantum mechanics of the final state interaction in order to fully understand such a suppression factor.

The effects on the particle distributions of the suppression factors $|g|^2$ and $|h|^2$ for baryon fragmentation is further discussed in section 6. Just as for the application to quark-jets in ref. [5] the noticeable effects are mainly connected to $(s\bar{s})$ production and therefore to strange particle formation.

4. The final state in the baryonic fragmentation region.

In this section we will consider the structure of the color force-field which we expect for a baryon fragmentation event (cf fig. 8). The three valence flavors of the original baryon will due to the reaction mechanism in section 2 be distributed along the force field with the L-quark (L for "leading") at the forward end of the field, followed by the J-quark (J for "joining" or "junction") and the I-quark (I for "initial" or "interacting"). The force field will be stretched according to fig. 9 between the corresponding colors (in this case red for L, green for J and blue for I). Note that the color of the interaction region, including the target, corresponds to the

initial color of the I-quark. Therefore, between the J-quark and the L-quark the joint force-field will have antired-red character (the I- and J-quarks act together as an antitriplet - antired - in color) while to the left of J it will have antiblue-blue character (with J and L acting together as an antiblue color antitriplet). Then for a pair $(q\bar{q})$ produced between J and L the q-particle is dragged towards J and similarly for a $(q\bar{q})$ pair produced between J and I the q-particle is dragged also towards J, i.e. the field changes direction at J.

Evidently the energy distribution in this field is very similar to the field in a quark jet. We will assume that the linear properties of the force field are the most important for the production process and that the breakup via $q\bar{q}$ pair production will happen in the same way as for the quark jets in section 3. The difference is then only the piece containing the J-quark where the field changes direction and where we will obtain a baryon instead of a meson. The resulting breakup structure will be according to fig. 10, in case there is a break also in the field between J and L.

Using a generation process of the same kind as the one outlined in section 3 it is evidently the position of the J-quark along the field (in scaled units x_J) which will determine whether a situation as in fig. 10 occurs. In that case the first rank hadron is a meson containing the L-flavor. If the first break occurs instead behind the J-quark, then the first rank hadron becomes a baryon containing the original diquark JL.

The analogues of figs 4 and 5 are shown in figs 11 and 12. The treatment of the J-quark position in terms of a single scale invariant coordinate x_J means that it is treated as moving with the velocity of light in the "backward" direction along the color field. This is an evident idealisation but we note that orbits with finite velocities in the cms as indicated in fig. 11 correspond in fig. 12 to orbits very close to the backward light cone at high energies.

In this way we treat the original LJ diquark structure by means of a probability distribution to find the J-quark at x_J when the L-quark is at 1. This distribution is related to the original diquark wave function. From the resulting pion spectra we find that the distribution must vanish for $x_J \approx 0$ and $x_J \approx 1$. We have not tried to make a best fit to experimental data but we find that the simple normalised distribution

$$\frac{dP}{dx_J} = 6x_J(1-x_J) \equiv \rho(x_J) \quad (2)$$

works well. The final state distributions are rather insensitive to variations around this shape. For the I-quark we have for simplicity assumed that its orbit is parallel to J and that its scaled position variable x_I is evenly distributed between 0 and x_J .

As in ref. [5] we assume that vector and pseudoscalar mesons are produced with ratio 1:1. For the baryon, however, the quarks must be in a symmetric state with regard to flavor and spin. (We assume that only baryons in the ground state 56-multiplet

are produced.) A basic idea behind the jet model is that the probability distribution is determined by the density of possible final states. Thus we assume the production of each baryon type to be weighted by the probability to find three quarks in a symmetric state with the proper quantum numbers, if the new quarks are produced with random spin and flavor (except for the "usual" suppression of s-quarks).

We assume equal probabilities for the three valence flavors in the proton to correspond to the I-quark, i.e. the one left centrally. With the $\rho(x_j)$ -distribution given in eq. (2) about 50% of the produced baryons will contain both quarks of the original leading diquark. We assume this di-quark to be in the states $uu(S=1)$, $ud(I=S=0)$, and $ud(I=S=1)$ with the probabilities $1/3$, $1/2$, and $1/6$ respectively, as given by SU(6), and we assume that if it is not split, its spin and isospin are left unchanged.

A $ud(I=S=0)$ di-quark is expected to be a more tightly bound system than the $I=S=1$ di-quarks. (This is connected to the $\Lambda^0 - \Sigma^0$ mass difference and probably to the difference in the u- and d-quark structure functions for $x_B \sim 1$ [10].) Thus if an interacting u-quark (I-quark) originally forms a $ud(I=S=0)$ with the d-quark in the proton, we expect this d-quark to be the J-quark, whereas the other u-quark, which is more loosely bound to the I-quark, will be the leading quark (L-quark) at the end of the force field. With SU(6) ratios for different di-quarks we thus find the probabilities $1/12$ and $11/12$ for the L-quark to be a d- or u-quark respectively.

The resulting distributions are not very sensitive to the precise value of this u/d-ratio for the L-quark except the π^+/π^- -ratio for z close to 1 (cf. the comment in section 6).

Noting the field structure between the L- and J-quarks in fig. 9 it is obvious that the suppression factor $|g|^2$ in eq. (1), which depends on the ratio of the available field length to the one necessary to produce the (transverse) mass of a $(q\bar{q})$ pair, should be taken into account in the field between J and L. Just as the effective field length, τ , in eq. (1) can be given a relativistically invariant meaning in the space-time difference between the origine and the production vertex (cf. fig. 6), a corresponding variable τ_J can be defined for the suppression factor $|g|^2((\kappa\tau_J)^2/\mu_\perp^2)$ relevant to the field between J and L. There is, however, an obvious question of principle involved here, due to our simple phenomenological treatment of the J-quark position by $\rho(x_J)$ in eq. (2). The field-line distribution may be more complicated around the J quark than in connection with the situation which in ref. [5] led us to the result of eq. (1). Fortunately the precise assumptions on the behavior of $|g|^2((\kappa\tau_J)^2/\mu_\perp^2)$ are not reflected in the final state particle distributions, although the general properties discussed in connection with eq. (1) are noticeable.

We will in the next section present the resulting model predictions both when we use the further suppression factor $|h|^2$ in eq. (1) and when we do not (the factor $|h|^2$ is independent of the change in field direction at the J-quark).

5. Comparison to experimental data.

In this section we present a set of comparisons between the model predictions for low $p_{\perp}$ proton fragmentation and the available experimental data.

In ref. [5] we have (in accordance with the observation in connection with eq. (1) above) found good agreement between the experimental data and the model for quark jets with a suppression factor $1/3 \sim 1/4$ for the strange relative to the nonstrange quarks. In the subsequent model calculations we have used the value $1/4$ for this factor.

For the model predictions we present two different curves, one where the factor $|h|^2$ in eq. (1) is equal to 1 (i.e. no suppression) and one where the expression

$$|h|^2(\alpha) = \frac{\alpha}{1+\alpha} \quad (3)$$

is used.

As the model is meant to describe only the particle distributions in inelastic nondiffractive events we have subsequently normalised all theoretical distributions to the cross section $\sigma_{\text{non-diffr}} = 26 \text{ mb}$ for pp collisions.

Figures 13 a-d show the longitudinal distributions for the inclusive production of $\pi^{\pm}$ and $K^{\pm}$. Reference [11] provides data for $\pi^{\pm}$ -production, already integrated over $p_{\perp}^2$, while refs [12,13] provide fits to the $p_{\perp}$ -distributions, which we have integrated. For the latter we show the statistical errors in the integrated points, calculated from errors in the fits. The

general behavior of the data is evidently well reproduced by the model and also the ratios between the different particles.

In fig. 14 we show as an example the ratio K^+/π^+ together with the data from refs. [12,14]

Baryon production is presented in figs. 15 a-c. In fig. 15a we give the invariant distribution $\frac{E}{\pi} \frac{d\sigma}{dp_L}$ for protons. Data is taken from ISR and FNAL refs. [15,12,16] and integrated over $p_{\perp}^2$ (the π^-p -data is scaled up by the factor $\frac{\sigma_{in}(pp)}{\sigma_{in}(\pi p)} \approx 1.5$)

The agreement is reasonable especially remembering that data includes diffraction and central ($p\bar{p}$)-production which is not accounted for by our model.

For $p \rightarrow nX$ there does not exist any high-energy, i.e. $\sqrt{s} \gtrsim 10$ GeV, data. However, assuming that $p \rightarrow nX$ is the same as $n \rightarrow pX$, we can use data on $pn \rightarrow p_{slow}X$ from FNAL, ref. [17]. We compare with these in fig. 15b. Here an extra uncertainty comes from the Fermi motion of the neutron in the deuterium target, which shows up in that the x-distribution extends beyond $x=1$. We also give the highest available $pp \rightarrow nX$ data, ref. [18] although the cms-energy is too low to make a comparison with our model really meaningful.

In fig. 15c we give the Λ -distribution compared to two sets of data, one from FNAL [19] and one from ISR [20].

Using our Monte Carlo generation program it is also possible to obtain predictions for any kind of correlations. Some comparisons with data from K^+p collisions at 32 GeV/c [21] are presented in figs. 16a-c. The π -production associated with a Λ -particle is considered in fig. 16a, where we give the π^+/π^- ratio for a Λ -trigger with $x \geq 0.4$.

In the data from ref. [21] protons with $p_{lab} > 1.2$ GeV/c are not identified but treated as pions. The pion spectra also contain a contamination from K^+ and K^- . In connection with a Λ -trigger the protons do not contribute because baryon-antibaryon pair production is neglected in our approach. However, for correlations with pion triggers the proton contributions cannot be neglected. In figs. 16 b,c we have therefore included the contributions from misidentified protons and kaons when we compare the experimental data with the predictions from our model (the full lines). The dashed lines show the corresponding results for the pion ratios if all particles could be properly identified. In fig. 16 b we give the ratio of the π^+ -distribution to the π^- -distribution for a π^+ or π^- trigger lying in the interval $.2 < x_1 = x_{trigger} < .4$. The variable $\tilde{x}_2$ is given by $\tilde{x}_2 = x_2/(1 - |x_1|)$, i.e. the ratio of the $\pi^\pm$ energy momentum to the maximum available value. The ratio of the $\pi^\pm$ -distributions for a π^- and a π^+ trigger is shown in fig. 16 c with similar notation. The model reproduces the qualitative features of the data.

We note in our model a strong correlation between the charges in the mesons and in the baryons, and that it is very important

to include unidentified kaons and protons to get agreement with the experimental data. We thus predict a very different correlation if the pions could be really identified, and we hope that the next generation of experiments will be able to quantify this difference.

Summarising the comparisons between our model and experiment in figs. 13-16 , we note that the model well reproduces the features of many different pieces of experimental information. The single particle inclusive spectra agrees within experimental errors with the model and also the features of the correlation data of ref. [21] are reproduced. We also note that the main features are very insensitive to the factor $|h|^2$ which is only important for the K^- -spectrum at large z -values. We will discuss this point further in the next section.

6. Concluding remarks

We have in this paper presented a possible reaction mechanism for low $p_{\perp}$ hadronic collisions and also a model for proton fragmentation in which the proton remnants stretch a linear color force field (or color flux tube) of much the same nature as in a quark jet.

The long flux tube, in which most of the particles are produced far away from the collision point is consistent with the absence of an intranuclear cascade in hadron-nucleus collisions.

It is also consistent with the observation that if all produced particles (baryons and mesons) are included, the z -distribution is very similar for proton and meson fragmentation.

The model contains a phenomenological function $\rho(x_j)$ which describes the structure of the leading diquark. This implies a basic difference between this model and a model of cascade character like the one described in ref.[22], where the diquark is treated as a structureless entity. In particular, if a first rank meson, containing the L-quark, is produced in our model, the structure of the remaining diquark is changed (i.e. the J-quark is then in general "closer" to the new end quark than to the original L-quark).

The suppression of the meson spectra relative to the baryon spectra for large z -values (like a factor $(1-z)^2$) is in this model due to the large probability that a hadron with large z also contains the J-quark and in that way becomes a baryon, irrespective of rank.

We have neglected the production of baryon-antibaryon pairs. We expect this to take place in the linear force field in the same way as in e^+e^- annihilation, and thus to be suppressed in the same way as in that case.

In our model the baryon always contains at least one of the original proton valence quarks, the J-quark. Thus e.g. the production of Ω 's should be suppressed, like the antibaryons are. The experimental cross section is also very small[23].

For a leptonproduction event we expect a very similar situation, with the difference that (for large enough x_B) the I-quark will now be the one struck by the photon or weak current, and will thus be at the current fragmentation end of the force field. However, because this is a hard process, it is possible that the struck quark cannot feel that it was more tightly bound to one of the remaining quarks than to the other. If so, in case of a struck u-quark ($\bar{\nu}$ -scattering and most e- or μ -events), the remaining ud diquark should be more symmetric in flavor, thus giving equal amounts of π^+ and π^- with large (negative) x_F .

Naturally our treatment of the force field as a purely longitudinal field is an oversimplification. Although we in this way obtain a very good overall agreement with data, we must expect corrections which can be important if we look in certain corners of phase space.

We have in particular completely neglected the so-called primordial transverse motion of the constituents inside the proton. In string language a description of the final state force field will then not be in terms of a purely one-dimensional string but instead a somewhat "bent" string system like the one described in ref. [25] for the target fragmentation region in a leptonproduction event. If this is the case there may be modifications in particular for the production of $(q\bar{q})$ pairs between the J- and L-quark. We note that the factor $|h|^2$ in eqs. (1) and (3) is most noticeable in connection with $(s\bar{s})$ pair production in this region and leads to a suppression of K^- mesons with large z . We have in ref. [5] discussed a few observable effects for strange particle production that will verify the $|h|^2$ -suppression hypothesis given in section 2. This discussion is applicable also to the baryon fragmentation model. Nevertheless we expect that the above-mentioned effect of the primordial transverse motion will be very similar to the $|h|^2$ -suppression and it will probably be difficult to disentangle one from the other.

The π^+/π^- ratio for large z -values is sensitive to the same details of the model. The experimental data imply the ratio $\pi^+/\pi^- \approx 5$ for $z \approx 1$ [12,14]. Of course our result depends on the treatment in section 4 of the more tightly bound ud ($I=S=0$) diquark, which gave the ratio $u/d = 11:1$ for the L-quark. We further note that in principle the function $\rho(x_j)$ could depend on whether the leading diquark is a ud ($I=S=0$) or a more loosely bound $I=S=1$ diquark. Because the favored meson spectra for $z \approx 1$

are correlated to $\rho(x_j)$ for $x_j \approx 0$ this could have an effect on the π^+/π^- ratio for $z \approx 1$. In table I we present results for the π^+/π^- ratio in the z -bin $0.8 < z < 0.9$ for various assumptions and we find that the differences are in general small.

We note that in this model the π^+/π^- ratio is not in any simple way related to the well-known Ochs-observation [24] that π^+/π^- as a function of z is similar to the structure function ratio in a proton u/d as a function of the Bjorken scaling variable x_B . The smoothly varying curve we obtain from a central value of around 1 to a forward value of 4-6 is in this approach a combined effect of the u/d ratio for the L-quark and the ratio of favored to unfavored particles in a jet.

Evidently an assumption like the one we have used that a ud ($I=S=0$) diquark is a more tightly bound entity than the corresponding $I=S=1$ diquarks is necessary both to understand the increasing (u/d) ratio as a function of x_B and the π^+/π^- ratio as a function of z .

We have investigated the influence of varying the distribution $\rho(x_j)$. In particular we have studied parametrizations like

$$\rho(x_j) = [Ax_j + B] \rho_0(x_j) \quad (4)$$

where $\rho_0(x_j)$ is given by eq. (2). Except for parameter values close to $B=0$ or $B/A = -1$ we find only very small differences in the particle spectra.

The presented results are also rather insensitive to our assumption that the parameter x_I , describing the I-quark,

is evenly distributed between 0 and x_j . Hopefully further studies of particle correlations can tell us more about what happens to the I-quark.

Concerning the spin of the produced quarks we have in the same way as for a quark jet [5] assumed that vector and pseudoscalar mesons are produced in the ratio 1:1 as indicated by experiments, rather than 3:1 as obtained from counting the available number of spin states. The results presented in section 5 do not include a possible similar suppression of the baryon decuplet relative to the octet. The inclusion of such a suppression factor would not significantly change the result for the stable baryons, p , n , and Λ or the meson spectra.

In conclusion we find that the resulting particle spectra are remarkably stable against variations of the details in the model, and that most of the properties are determined by the basic assumption about a linear force field of the same nature as in a quark jet.

We have in this paper only considered the distributions in longitudinal momentum and energy, and we will in forthcoming publications come back to the transverse momentum and spin properties.

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Figure captions

1. The color of a quark is, due to asymptotic freedom, not well localized, but rather distributed like a blob around the quark. A meson can be imagined as two blobs confined in a bag as shown in the figure. It is thus e.g. red in one end, antired in the other, and "white" in the middle.
2. Consider two meson bags approaching each other (the bags are Lorentz contracted). If the colored parts overlap, the two bags can unite to form one larger bag, which is stretched to a color flux tube. The bag pressure retards the meson remnants. The colliding blob will lose all its momentum before the other blob will be retarded. The bag is thus stretched to the same length as if all the energy were carried by one of the valence quarks.
3. If a quark and an antiquark go out from one point with large momenta in opposite directions, a color force field is stretched between them. This field is broken by the production of quark-antiquark pairs with zero energy, which are pulled apart by the field.
4. The final result with asymptotically emerging $q\bar{q}$ states (mesons).
5. The situation in a Lorentz frame which is boosted along the leading quark. The original quark, q_0 , starts with fairly low momentum and will soon be stopped by the force field, whereafter it turns around and proceeds in the same direction as $\bar{q}_0$. It combines with the antiquark $\bar{q}_1$ to form the first rank meson. The Feynman z -variable for this meson is determined by the distance z_1 .
6. The production of a heavy mass ($\mu_\perp$) $q\bar{q}$ pair in the force field from the original $q_0\bar{q}_0$ pair which starts out in the space-time point O . The invariant space-time distance τ from O to the production vertex V is smaller than $\mu_\perp/\kappa$ and the $q\bar{q}$ pair cannot classically be produced at rest inside the field.
7. The production of a heavy mass ($\mu_\perp$) $q\bar{q}$ pair such that the meson ($\bar{q}_2q$) is much slower than the next-rank meson ($\bar{q}q_1$). The invariant space-time distance τ' between the production vertex V' and the classical endpoint O' of the field obtained by the production of the $q_2\bar{q}_2$ pair at the vertex V_2 and the $q_1\bar{q}_1$ pair at the vertex V_1 is small compared to $\mu_\perp/\kappa$.
8. The analogue of fig. 2 when one of the particles is a baryon.

9. The force field between the separated colors. The blue and green quarks act together as an antitriplet-antired-in color, while the green and red quarks act as an antiblue color antitriplet.
10. The resulting breakup structure in the baryon fragmentation region.
11. The analogue of fig. 4 for baryon fragmentation. The trajectory of the J-quark is indicated. In this case it happens to go into the second rank hadron, which thus becomes a baryon.
12. The analogue of fig. 5 for baryon fragmentation. In this Lorentz frame the J-quark moves "backwards" with very large velocity, i.e. close to the light cone.
13. Inclusive $\pi^\pm$ and $K^\pm$ spectra in pp collisions
a-d — : Model prediction with suppression factor $|h|^2 = 1$
---- : - " - $|h|^2$ as given in eq.(3).
 - : Data from Kafka et al. [11] $\sqrt{s} = 19.7 \text{ GeV}$
 - : Data from Taylor et al. [13] $\sqrt{s} \gtrsim 10 \text{ GeV}$, integrated over $p_\perp^2$
 - : Data from Johnson et al. [12] 100-400 GeV/c, integrated over $p_\perp^2$
14. The ratio π^+/K^+ for inclusive spectra.
Curve notation as in fig. 13.
 - : Data from Johnson et al. [12], $p_\perp \leq 1.5 \text{ GeV}/c$
 - : Data from Singh et al. [14], $p_\perp = 0.65 \text{ GeV}/c$
15. Inclusive p, n, and Λ spectra in pp collisions.
Curve notation as in fig. 13
a) p : ○ Data from ISR [15] integrated over $p_\perp^2$
+ Higgins [16] π^-p 360 GeV/c
/// Data from Johnson et al. [12] with the $p_\perp$ -integration estimated.
b) n : ○ Eisenberg et al. [17] 195 GeV/c (actually pn → pX)
+ Blobel et al. [18] 24 GeV/c
c) Λ : ○ Kichimi et al.[19] 400 GeV/c
+ Erhan et al. [20] ISR.

16. Model predictions for two-particle correlations.

The dashed lines are true π -ratios, the full lines include misidentified kaons and protons.

- a) The π^+ to π^- ratio associated with a Λ -trigger. Data from De Wolf, e. a. [21].
- b) The π^+ to π^- ratio associated with a $\pi^\pm$ -trigger; both particles in the proton fragmentation region. Data from de Wolf, e.a. [21]
- c) The ratio of π^+ -production associated with a π^- -trigger to that associated with a π^+ -trigger. Data from De Wolf, e.a. [21].

Table I

Gives the variation of π^+/π^- with the u/d-ratio for the L-quark. Both $|g|^2$ and $|h|^2$ corrections (cf eq(1)) are included.

$(u/d)_L$	9:1	11:1	13:1	Data [14]
π^+/π^- $0.8 < z < 0.9$	4.9 ± 0.7	5.4 ± 0.7	5.8 ± 0.7	$4.5 - 6.0$



red



antired



"white"

Fig. 1

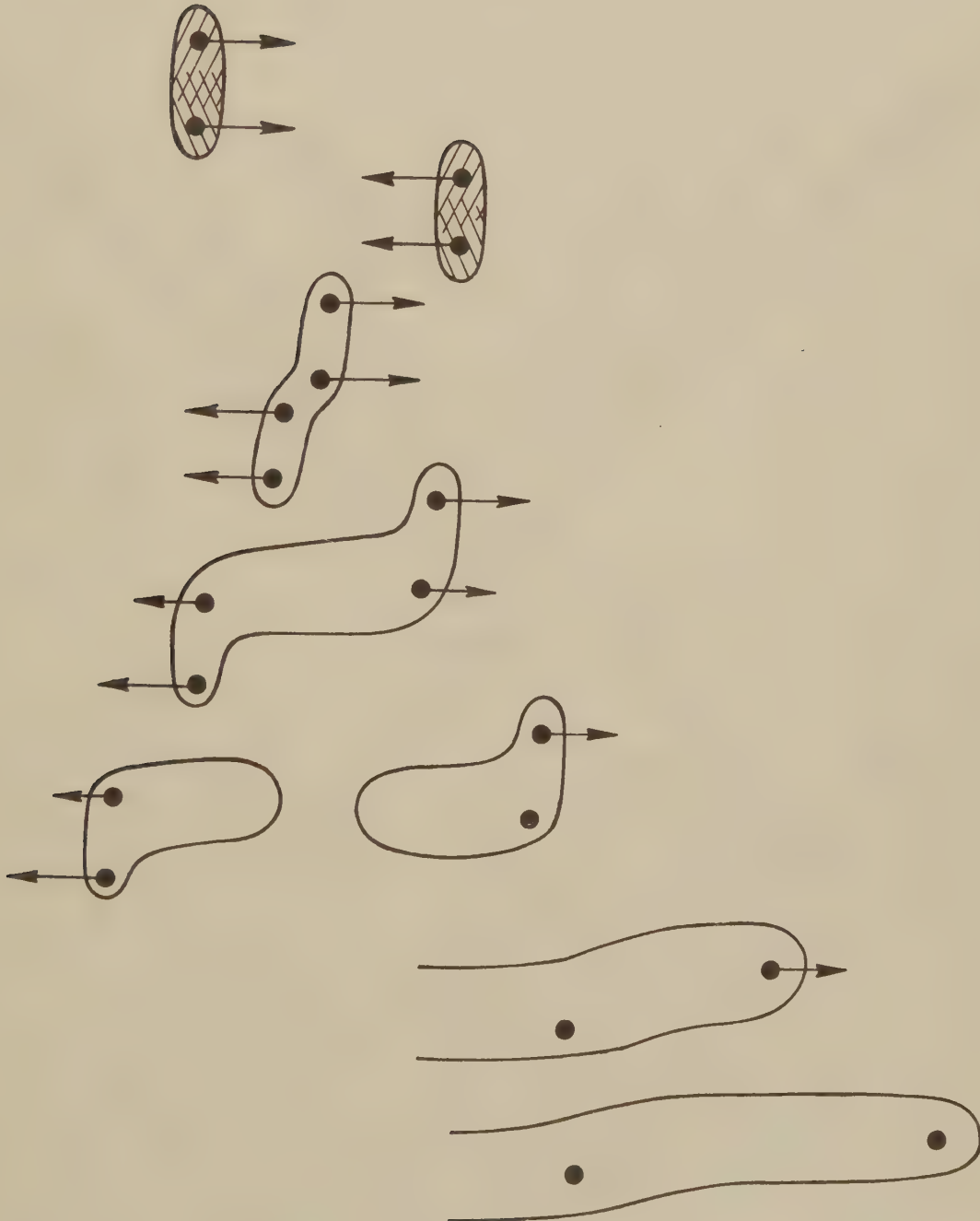


Fig. 2

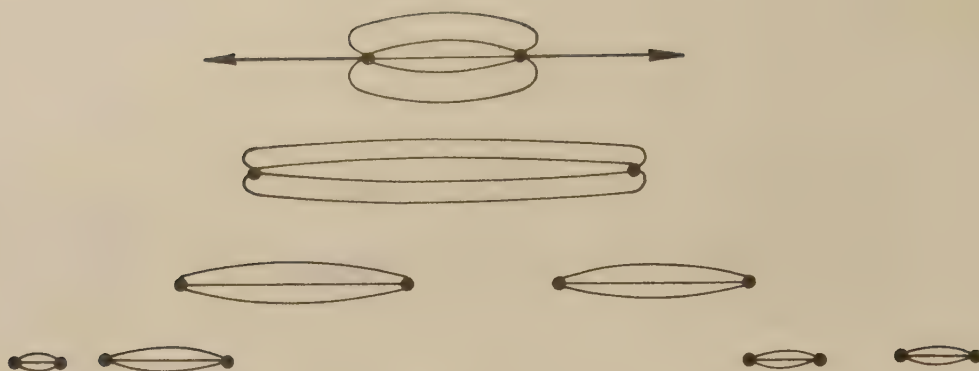


Fig. 3

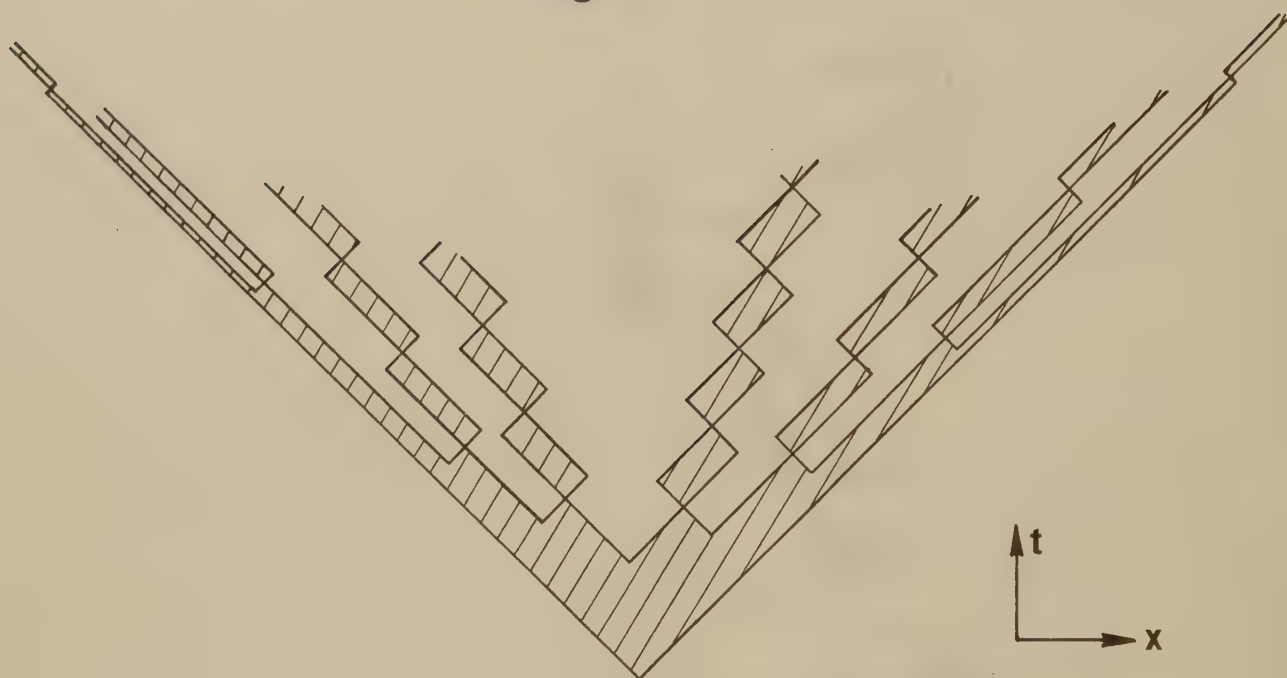


Fig. 4

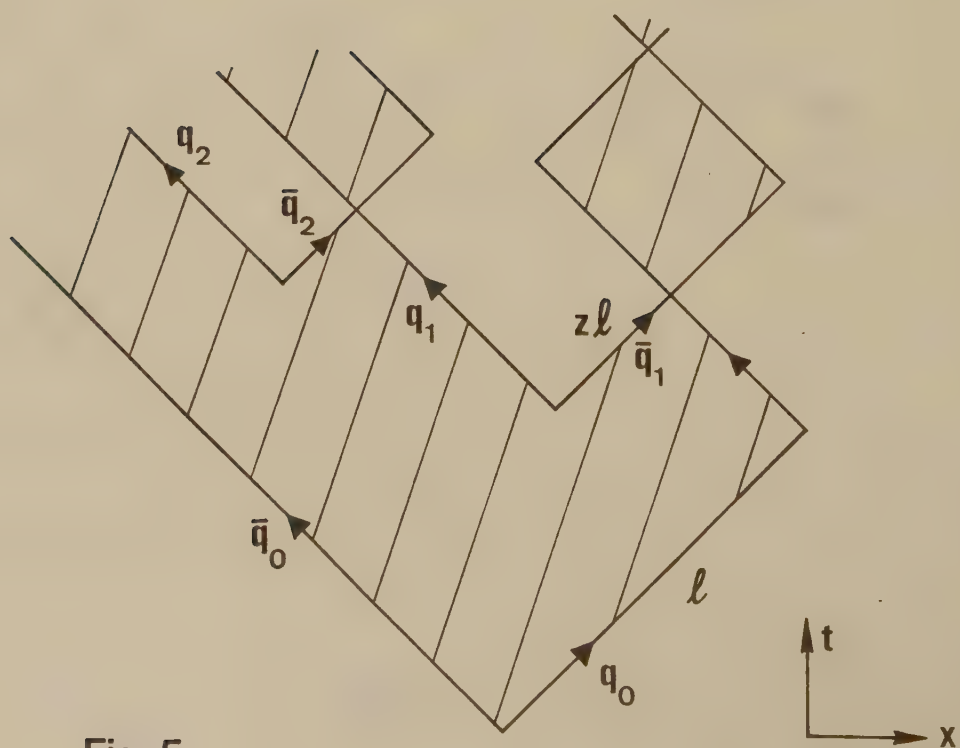


Fig. 5



Fig. 6

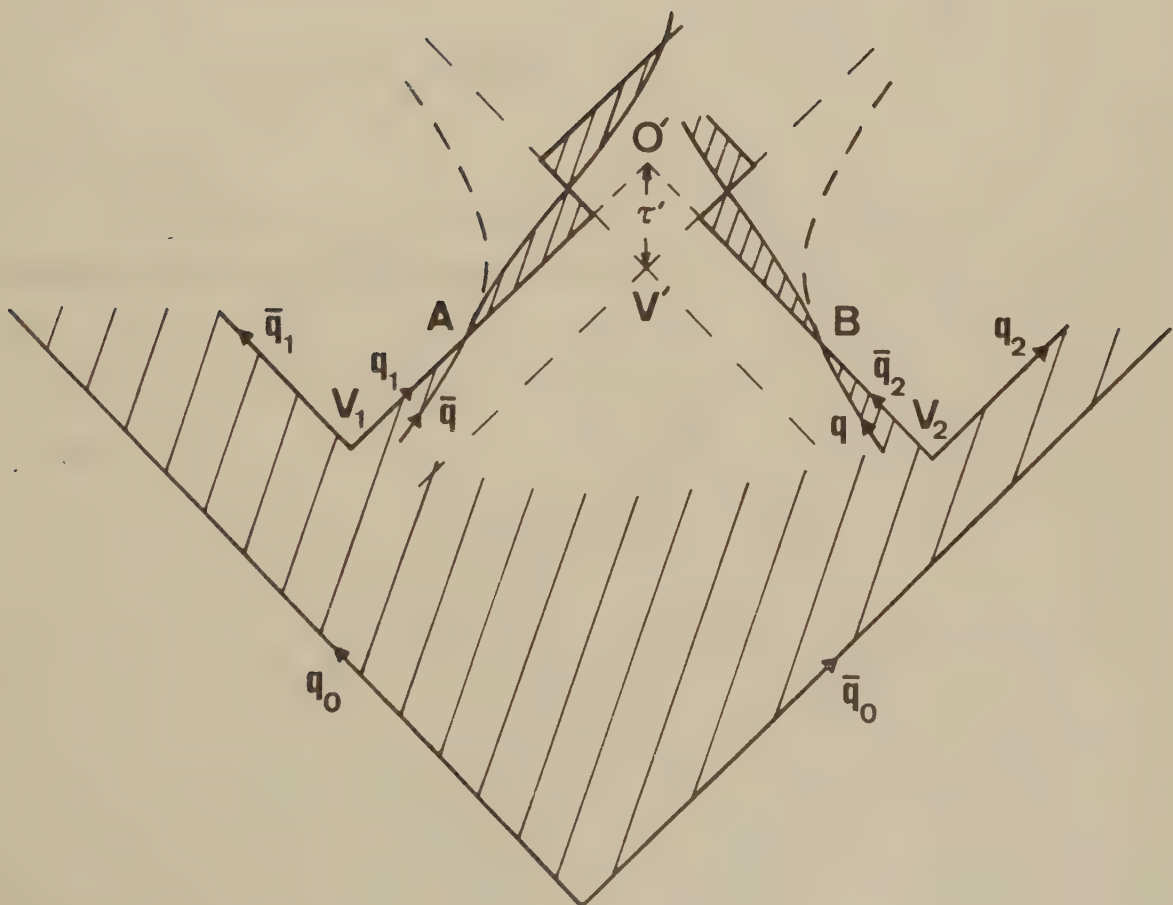


Fig. 7

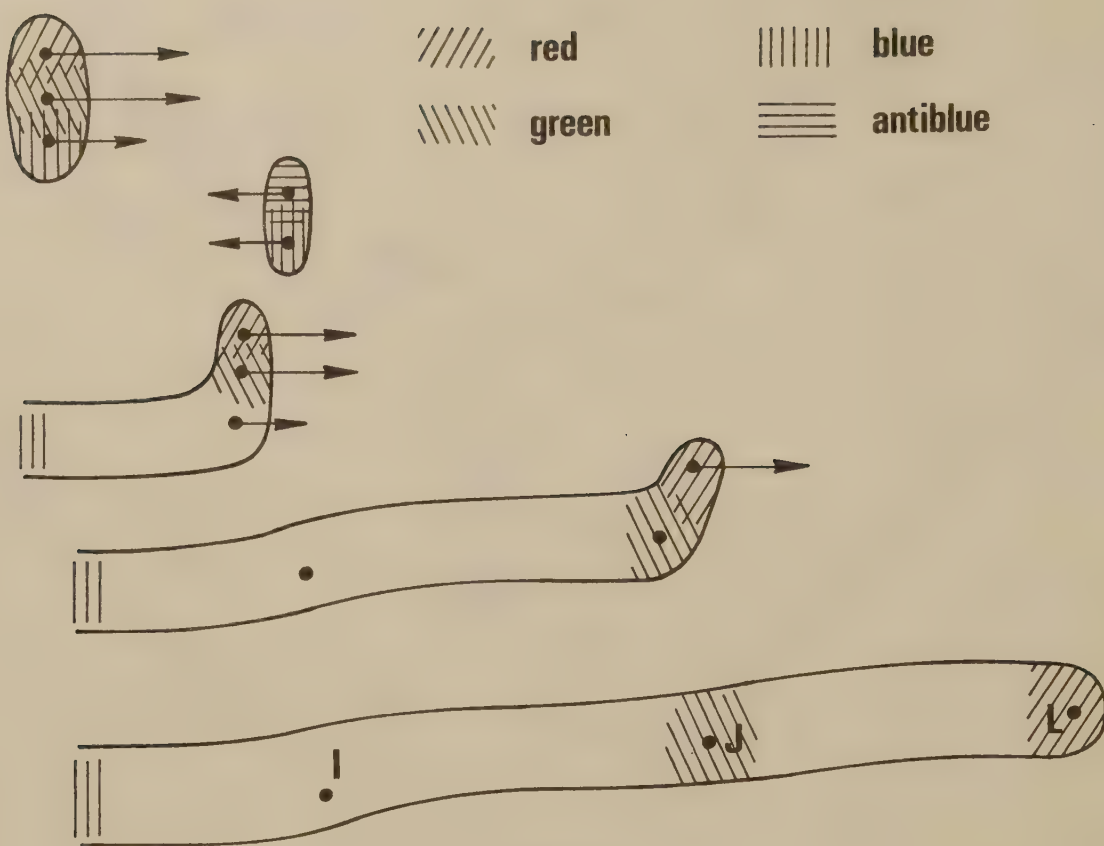


Fig. 8

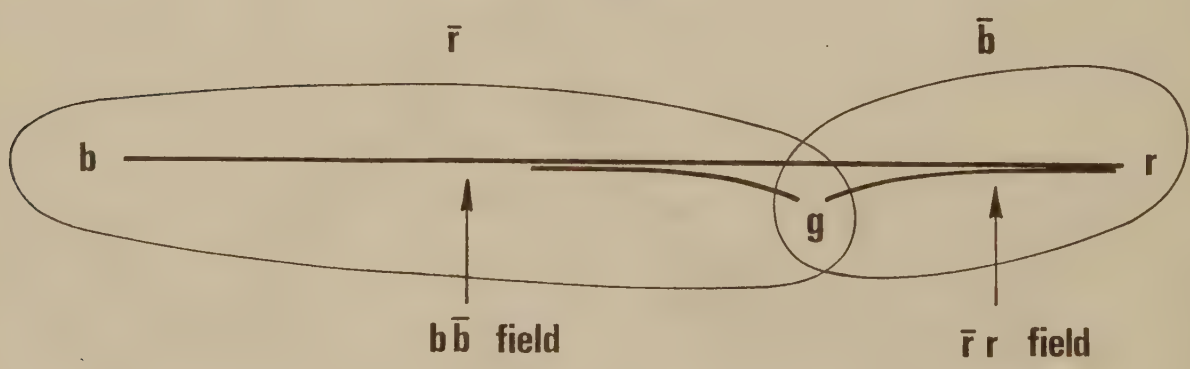


Fig. 9

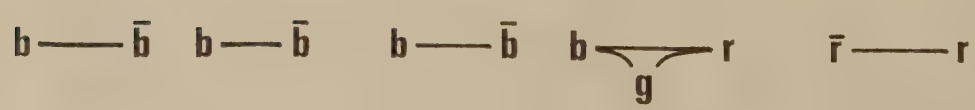


Fig. 10

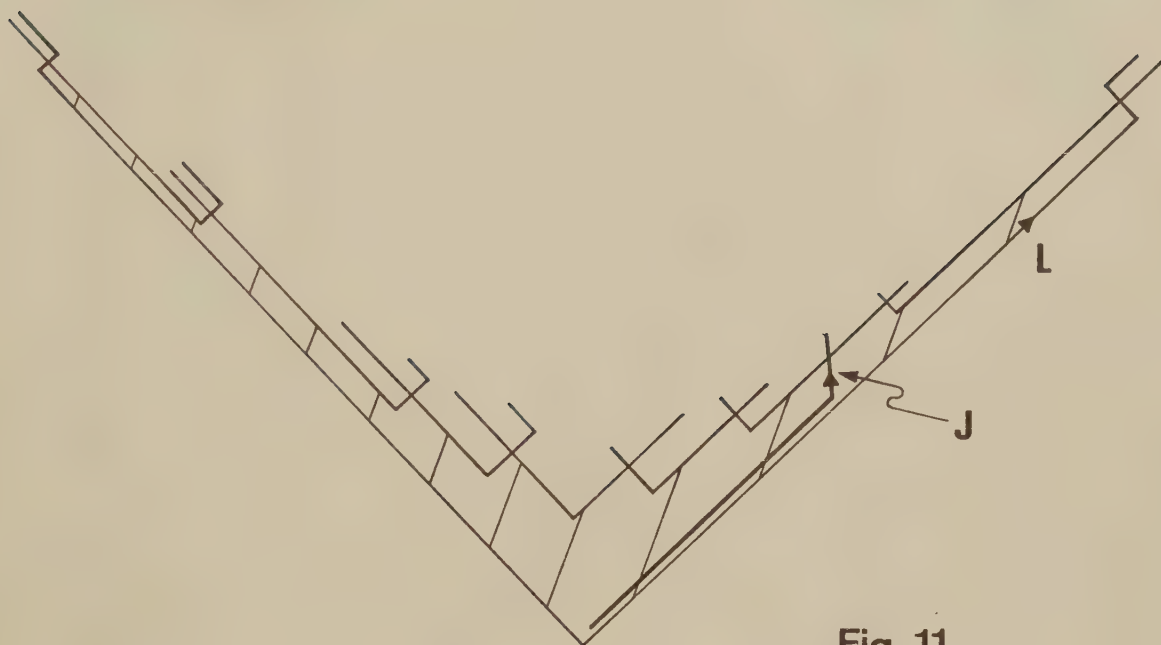


Fig. 11

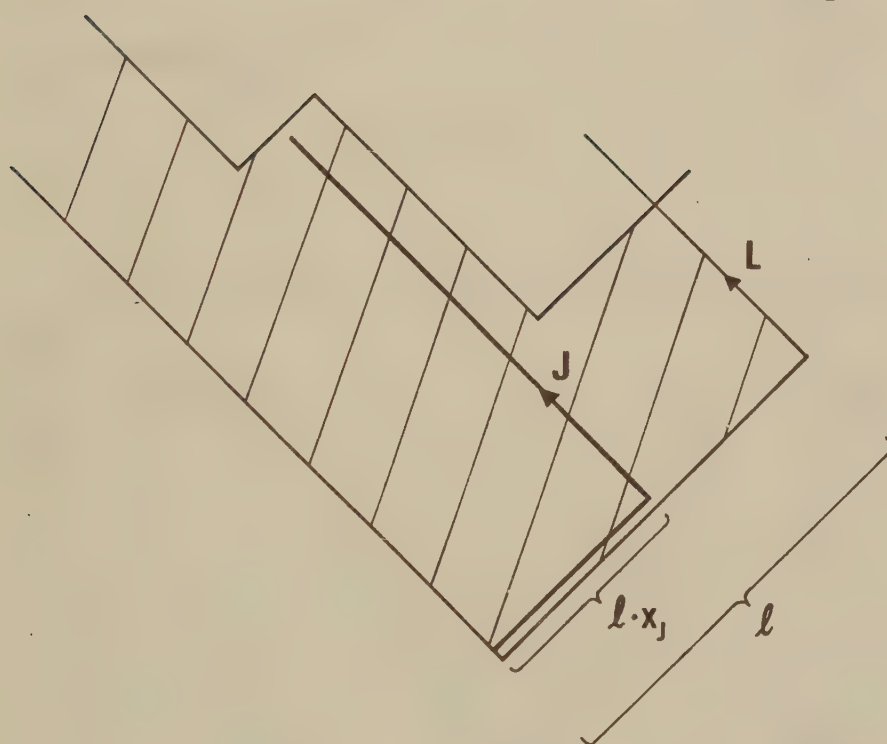


Fig. 12

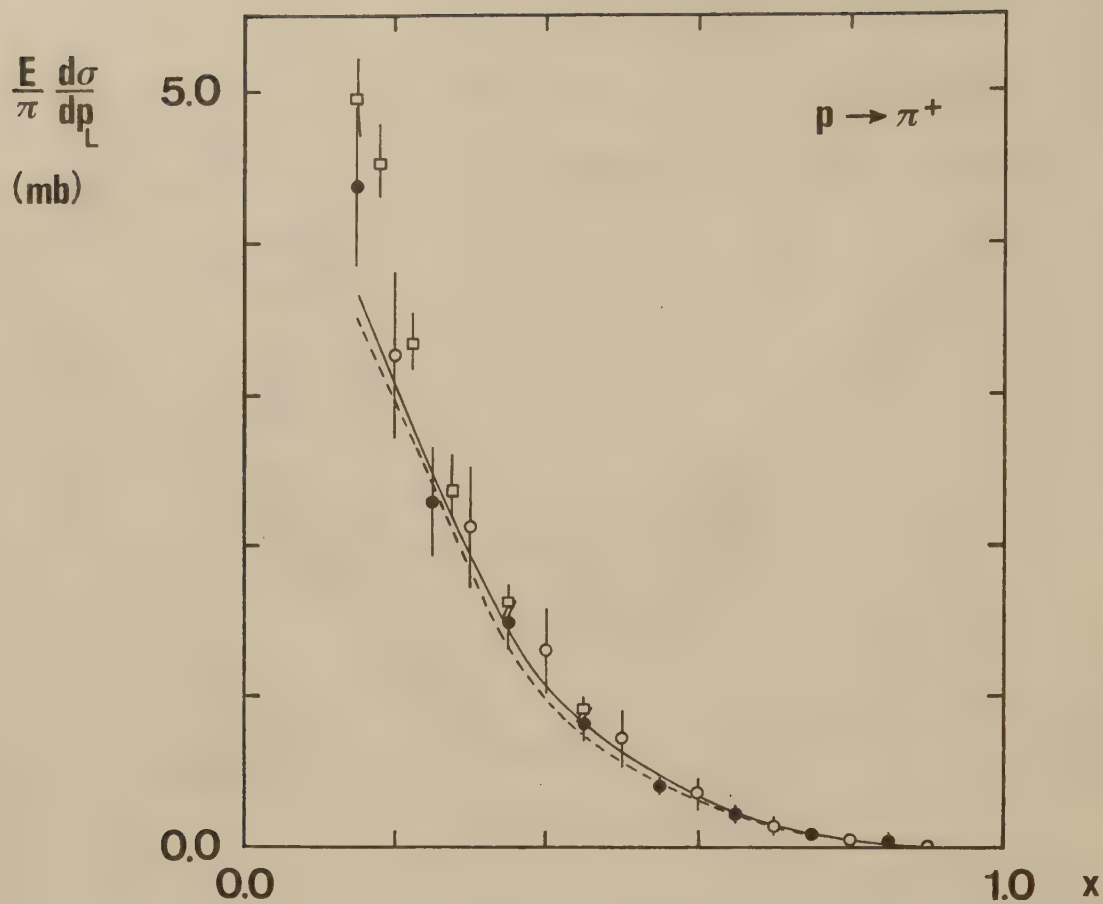


Fig. 13 a

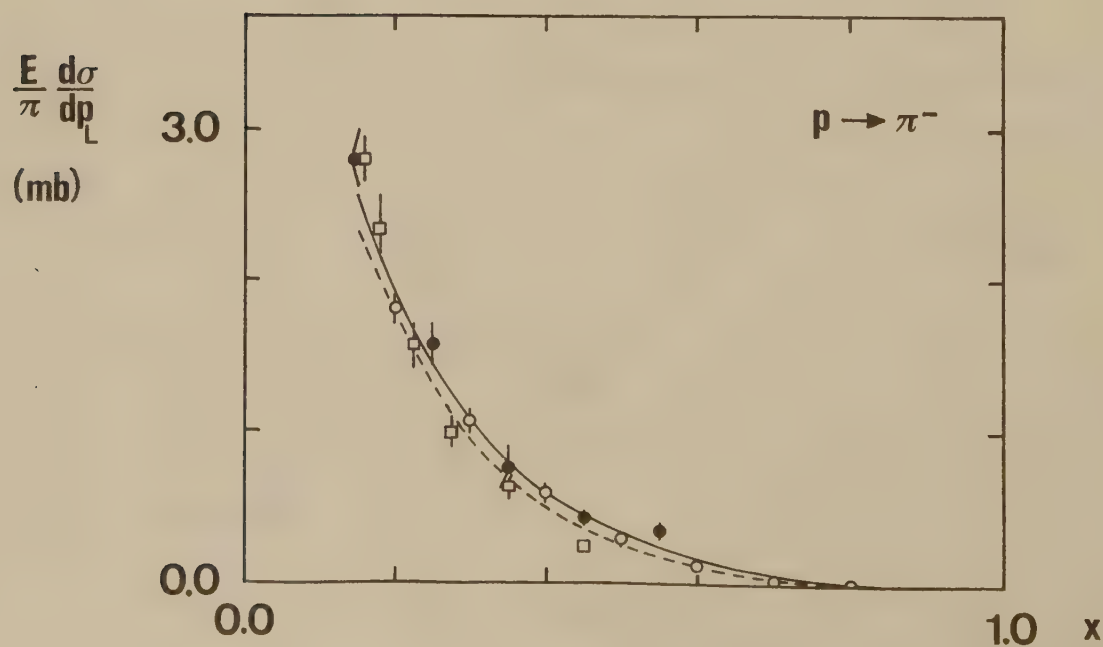


Fig. 13 b

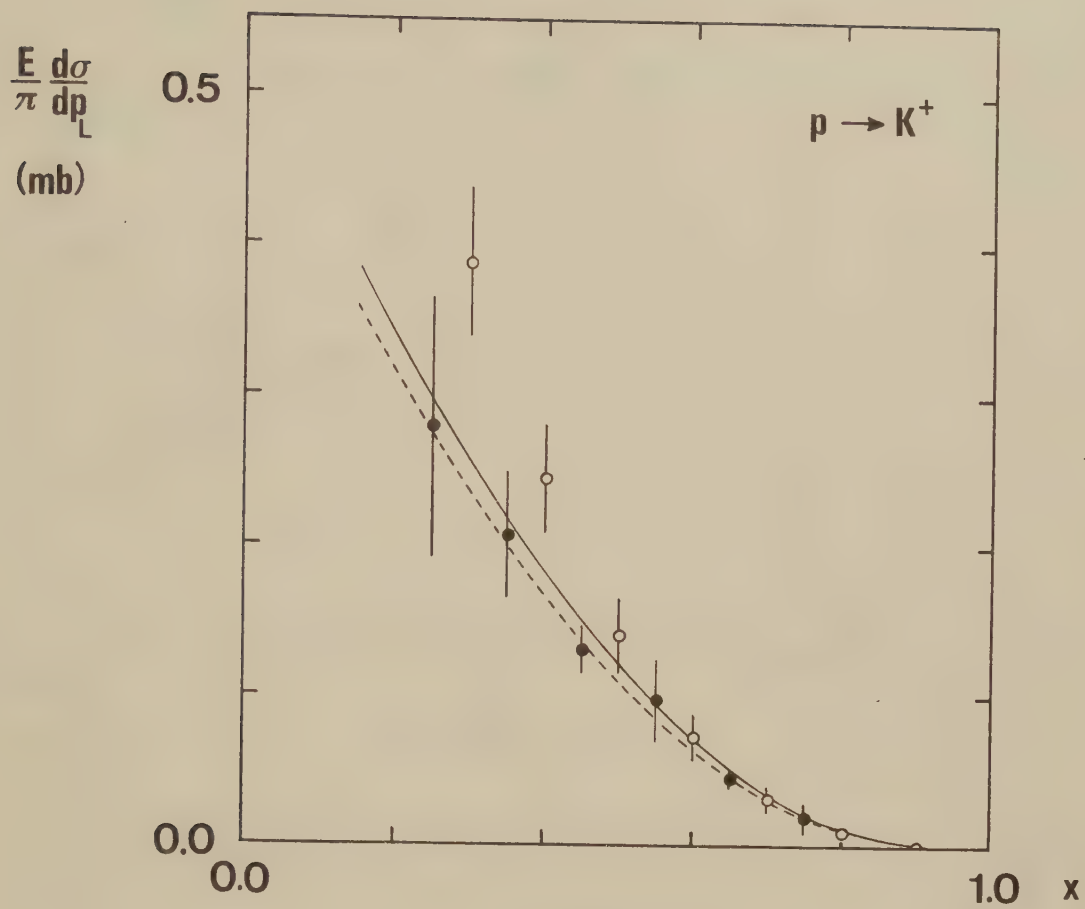


Fig. 13 c

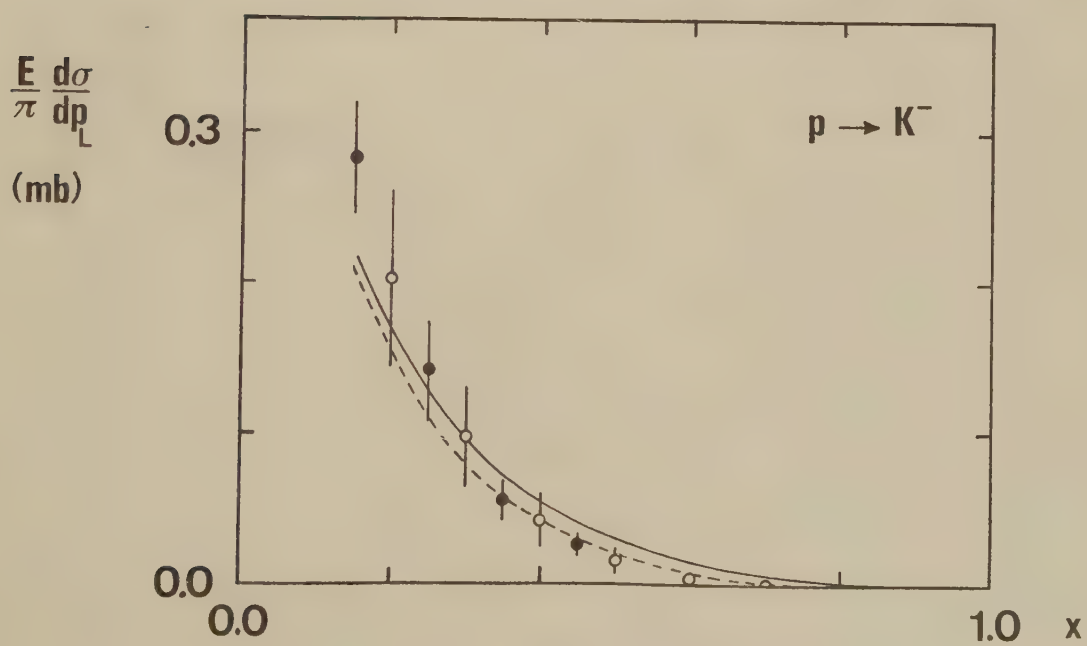


Fig. 13 d

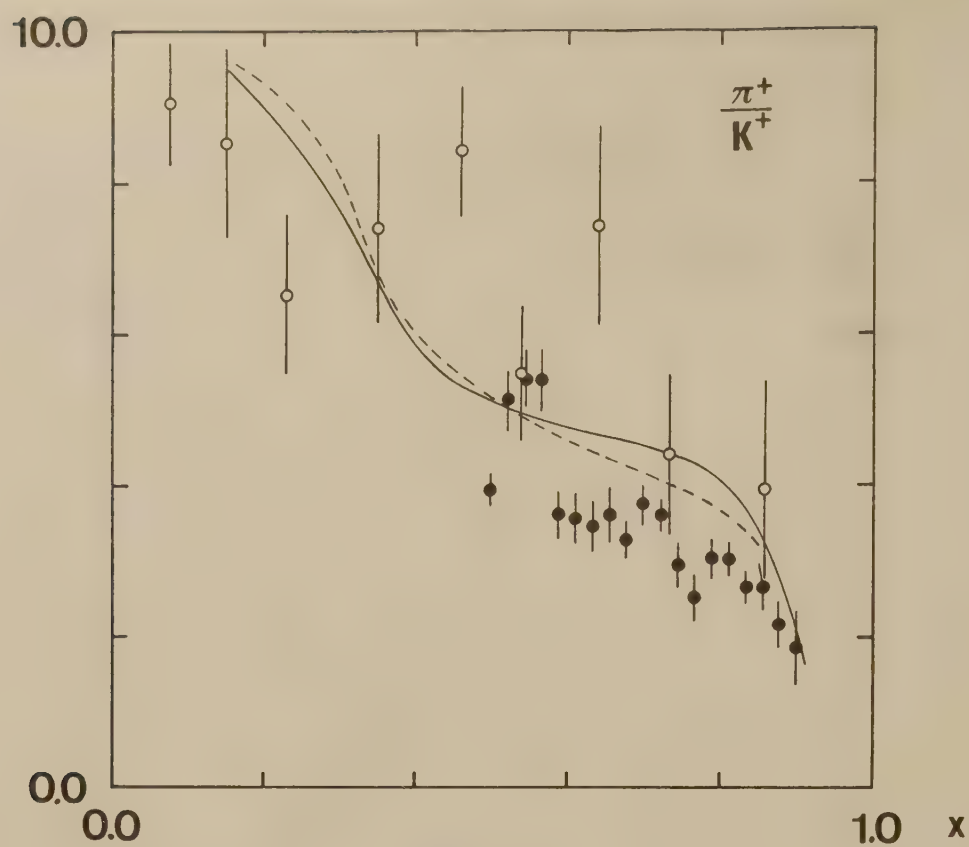


Fig. 14

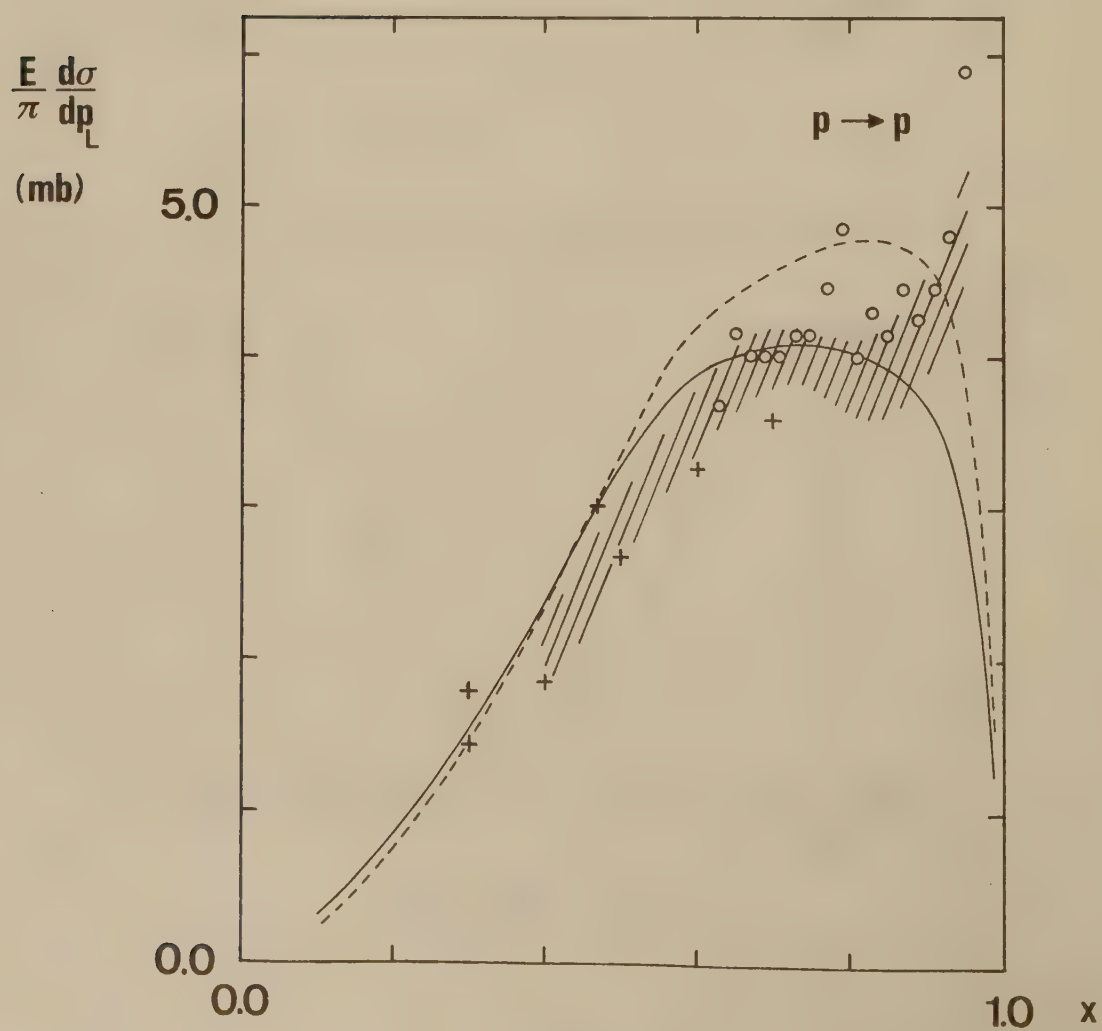


Fig. 15 a

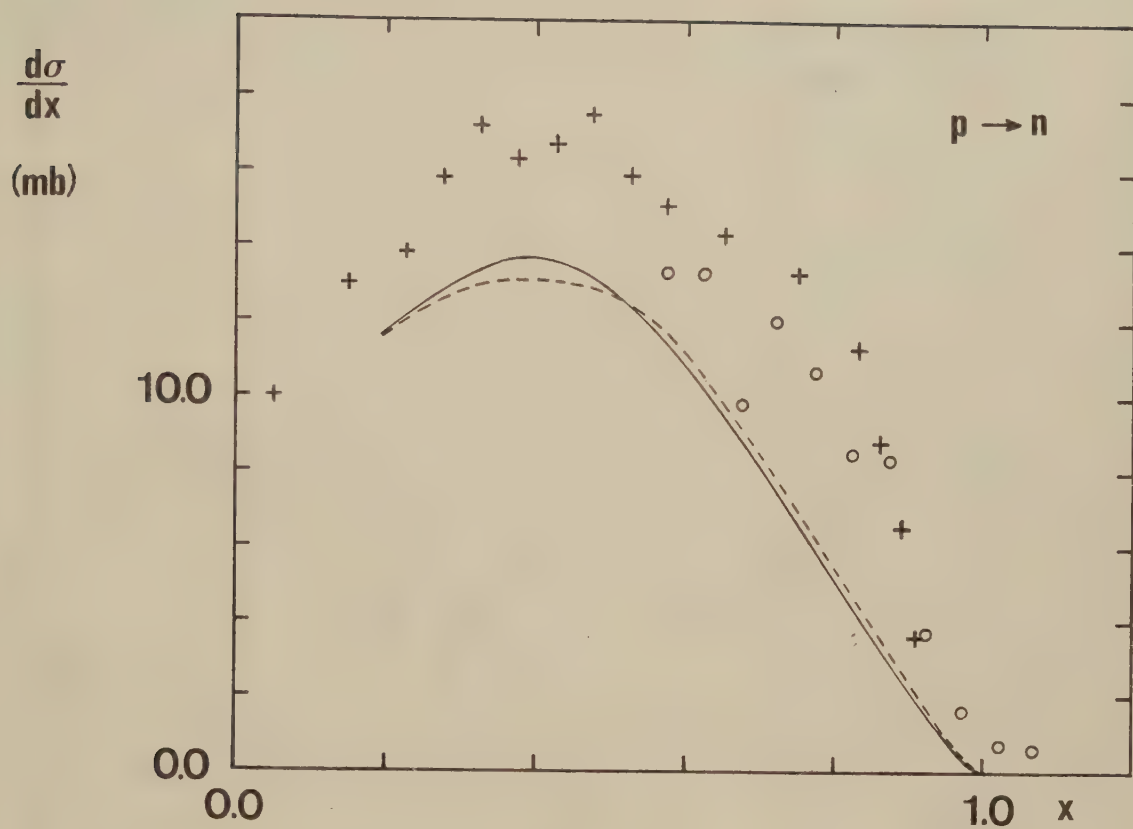


Fig. 15 b

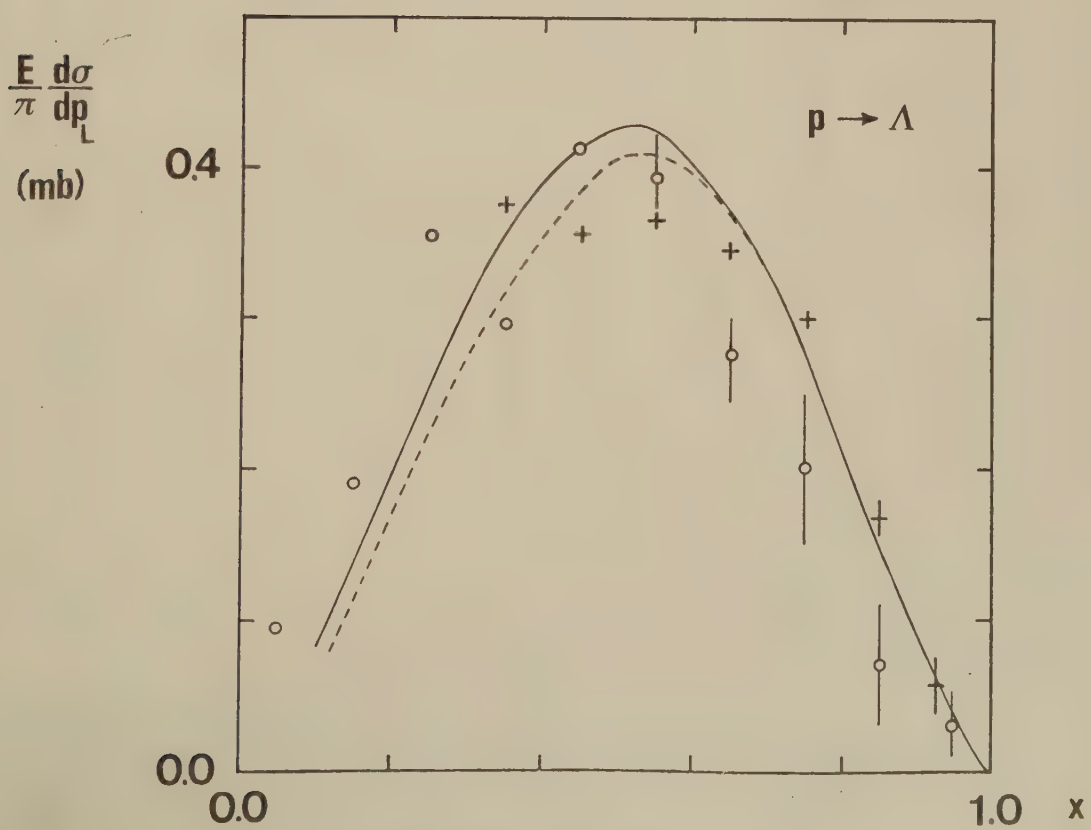


Fig. 15 c

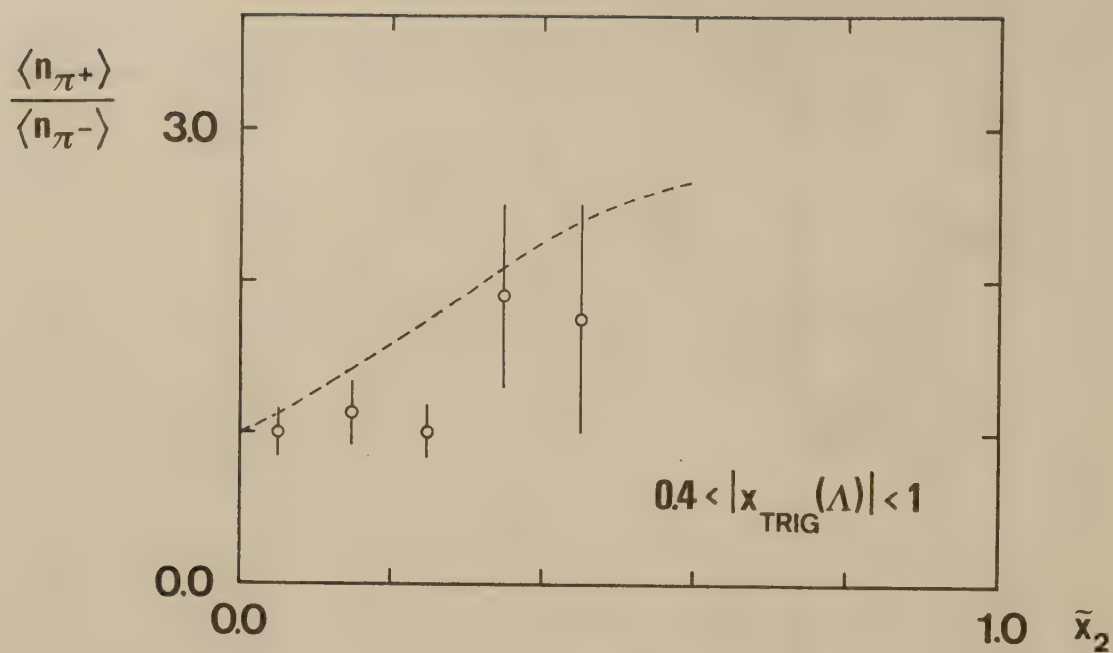


Fig. 16 a

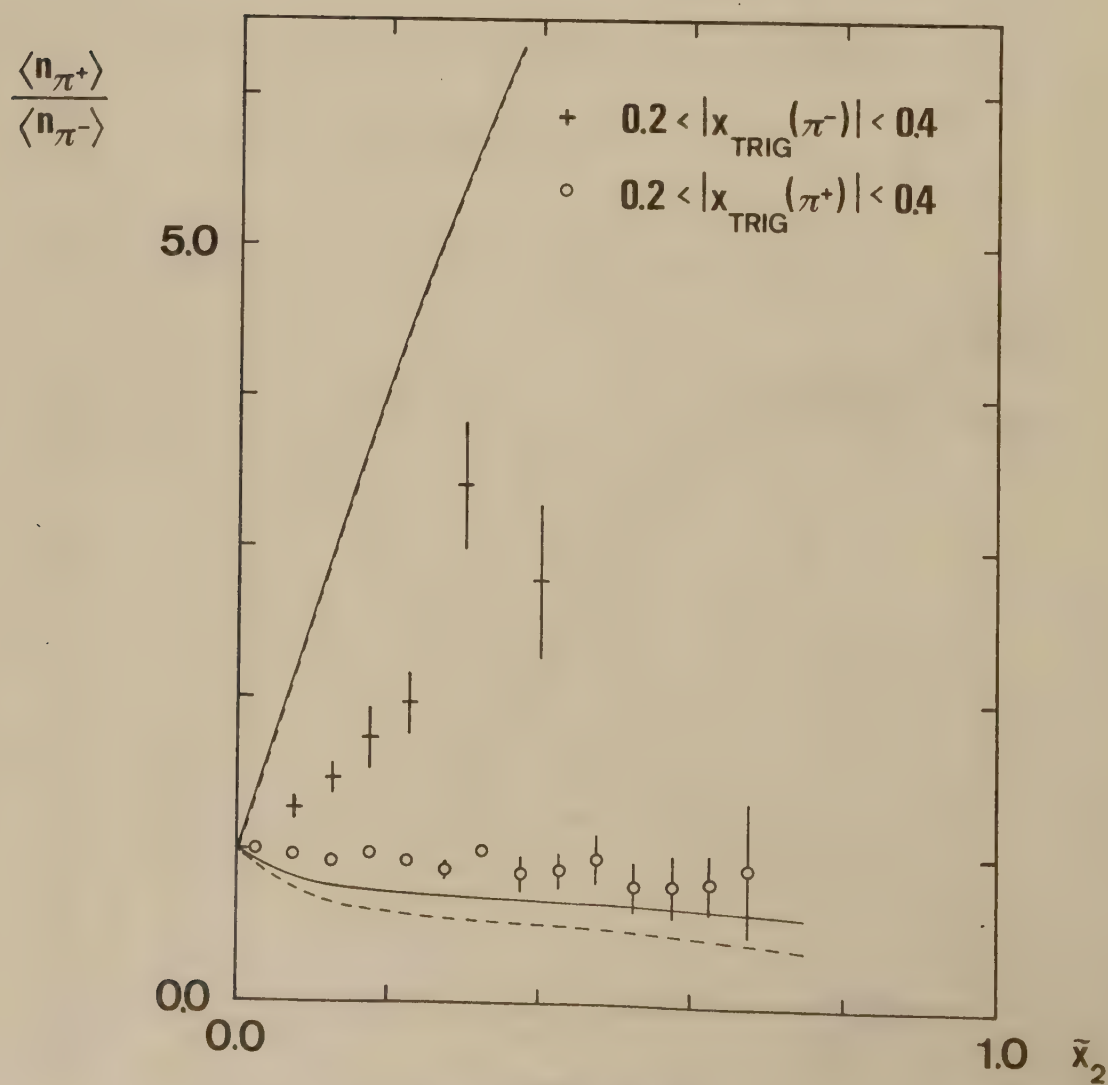


Fig. 16 b

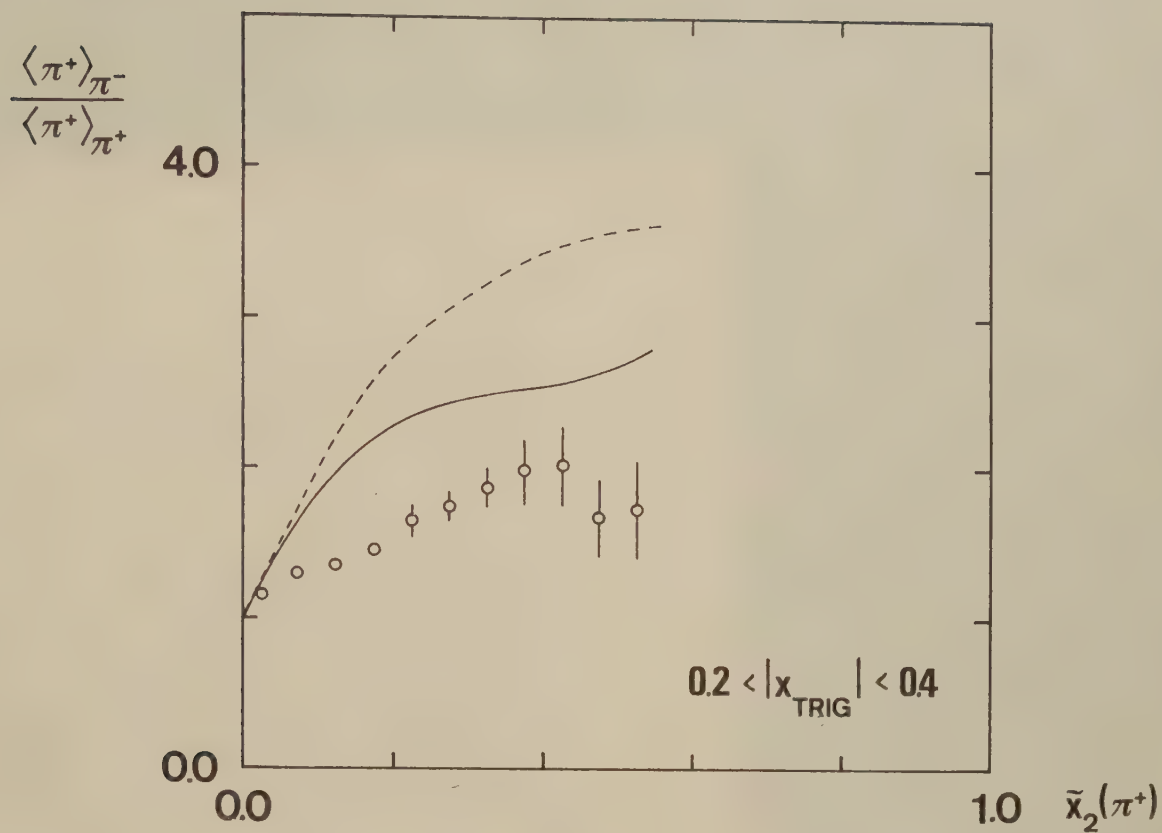


Fig. 16 c

in hydrodynamic models has decreased in deference to QCD-inspired quark fragmentation models also for processes involving low momentum transfers as in ordinary hadronic collisions. This trend is reflected in this thesis too where a model of the latter type is presented in paper III.

To understand these quark fragmentation models we must consider what happens when we try to liberate a quark from the rest of a hadron. Since quarks are glued together by a potential that grows linearly with distance, it would require an infinite amount of energy to separate them completely from each other. If we try to do this, it will ultimately become energetically favorable for the interaction energy to materialize in form of the creation of new quark-antiquark pairs. These will then, together with the original quarks, form new particles. Thus emerges a picture of the fragmenting quark: if a quark is given a large impulse in a direction different from the other quarks in a particle, it will not become free itself, but will appear as the creation of a number of particles in the approximate direction originally given to it, i.e. it emerges as a jet of particles in that direction.

These jets are seen in many types of high energy collisions: in deep inelastic lepton-hadron scattering, e^+e^- -annihilations and large- $p_{\perp}$ hadronic collisions, where quarks are assumed to be violently hit and kicked out of a hadron. But even in ordinary low- $p_{\perp}$ hadronic events there are, as mentioned earlier, two jets produced.

When two hadrons collide we can, as in paper III, think of the interaction as creating a connection, a color tunnel, between two of the quarks, one from each hadron. The connection makes these quarks lose their momentum while the other quarks of the hadrons continue undisturbed. This leads to a stepwise stretching-out of the bag to a tube or stringlike configuration, which then breaks down into small pieces as new quark-antiquark pairs are created. In this way the

outcoming low- $p_{\perp}$ jets are the results of a comparatively slow stretching-out of a color force field which then breaks down into pieces. This is to be compared to the more violent processes that produce the jets in deep inelastic scattering and e^+e^- -annihilation.

Other attempts to understand ordinary hadronic collisions in terms of quarks have also been made. For instance in the middle of the 70's Anisovich and Shekhter [3] used a naive additive quark model to obtain final particle distributions. Recently there has also been some interest in what is called recombination models. (The model proposed by van Hove and Pokorski (see below) [4] is one example of such a model.) Here the valence quarks, carrying the quantum numbers and most of the momentum of the particle, are thought not to take part in the interaction, but to fly through and then to recombine with some other quark that is around into a leading particle.

These models can explain some of the data in the fragmentation regions, but the advantage of the model presented in paper III is that it contains a plausible and complete reaction mechanism for hadronic collisions as well as explaining the experimental data.

One feature of high energy hadronic collisions that came as a surprise for most physicists is that Λ -particles produced with high momentum are polarized. This has a quite natural explanation in the model of paper III, but it is difficult to understand in most other models. In the future, polarization phenomena may well prove to be a main source for discriminating between different models and thus teach us more about the fundamental dynamics of hadronic collisions.

SUMMARY OF PAPERS

Some of the basic assumptions in the hydrodynamic model depend on what view of the underlying constituent dynamics we have. Therefore it is interesting to find out how stable the model is against variations of these assumptions.

In paper I we investigate variations of the initial state of the Landau model. There the colliding protons are thought to be stopped completely by the strong interaction. After the passage of shock waves, the thermal equilibrium mentioned earlier is established. The initial state for the solution of the hydrodynamic equations therefore is given by a uniform distribution of matter with the extension of two Lorentz-contracted nucleons. A different view of what happens during the collision has been given by v. Hove and Pokorski [4]. They picture the proton as consisting of three valence quarks carrying the quantum numbers and a neutral remainder called the glue. When then two protons collide, the quarks pass by each other and leave most of the glue behind interacting and in an excited state. The glue is thus distributed in the way that some of it is stopped completely and at high temperature, and some of it is dragged along by the valence quarks and has a lower temperature.

Assuming these kinds of initial distributions of the matter we solve the equations of motion and obtain a direct correspondence between the initial state and the final particle distribution. The investigation shows that the introduction of even fairly slow-moving parts in the initial distributions gives a considerably flatter final distribution. Thus the model is rather unstable against variations of the initial conditions.

Another crucial assumption in hydrodynamic models is the form of the equation of state for the fluid. In terms of the velocity of sound c_1 in the fluid it can be formulated as $c_1^2 = \frac{dp}{d\epsilon}$ where p = pressure and ϵ = energy density in the cms. Landau originally assumed a constant c_1 equal to $1/\sqrt{3}$ as for a free massless relativistic gas. Some later authors have tried to fit $c_1 = \text{const}$ to the experimental data. Paper II contains an investigation of the more general case when c_1 is allowed to vary with temperature T , i.e. $c_1(T)$. The generalised equations of motion are derived and solved in a special case. The solutions obtained thereby are then compared to free

relativistic gases, and it is found that for reasonable constraints on the physical properties of these gases, the T -dependence of c_1 can be neglected. We conclude that the LHD is stable to this kind of variation.

In paper III we present an elaborate model for the fragmentation of baryons. We imagine two hadron bags colliding. The interaction creates a connection, a color tunnel, between them, which makes the bag become stretched out and the valence quarks of the hadron lose their momenta, one at a time. For the baryon this results in a distribution of the valence quarks where one quark is fairly close to the interaction region, one at the other end of the stretched bag and the third somewhere between the other two. In the paper we give a detailed description of the model and describe a Monte Carlo program, which we have constructed, to generate jets in accordance with it. These jets give final particle distributions along the beam axis in good agreement with data, both for mesons and baryons. We also present some non-trivial correlation effects to be tested in future experiments.

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